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MICROWAVE EMISSION FROM HIGH-TEMPERATURE PLASMAS*

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Recent investigations[1,2] have suggested alarmingly high energy losses from plasmas contemplated in fusion research due to cyclotron radiation. Details will be published elsewhere of a calculation we have made of this effect in which it is concluded that the loss is definitely negligible for plasma temperatures less than and most probably greater than 100 kev. The emission for the fundamental frequency $(eH/m_e c)$ and the first few harmonics[3] is sufficient, however, to be detectable and holds promise of being a very useful plasma "thermometer" to measure the electron temperature.

If an electric field, $E = E_0 e^{i\omega t}$, (where ω is $> \omega_p$, the plasma frequency) is incident on a magnetically confined plasma, the equation of motion in the x direction for the electrons can be shown to result in the first-order nonrelativistic equation:

$$-m_e\omega^2\left[1+\tfrac{3}{2}\left(\frac{v}{c}\right)^2\right]\frac{P_{x\gamma}}{eN(v)S_\gamma(v)} = eE_x - ig\omega\frac{P_{x\gamma}}{eN(v)S_\gamma(v)} + \frac{im_e\omega\gamma\omega_0 P_{y\gamma}}{eN(v)S_\gamma(v)}, \qquad (1)$$

where m_e is the rest mass of the electron, (v/c) is its velocity relative to light velocity, e is its charge, ω is the frequency of the incident radiation, g is a collision damping term, $\omega_0 = (eH/m_e c)$ where H is the confining magnetic field, E_x is the amplitude of the incident wave, $N(v)$ is the Maxwell-Boltzmann distribution of the electron velocity, and $S_\gamma(v)$ is the coupling strength of the γth harmonic frequency to the electron relative to the fundamental. The plasma polarization for the γth harmonic is

$$P_{x\gamma}\,dv = S_\gamma(v)\,N(v)\,\xi(v)\,dv, \qquad (2)$$

where $\xi(v)$ is the electron displacement along the x axis for an electron of velocity v. The total polarization is

$$P_x = \int_0^\infty \sum_{\gamma=1}^{\infty} P_{x\gamma}\,dv. \qquad (3)$$

Substitution of Maxwell's equations into Eq. (1) results in an elementary second-order differential equation from which the index of refraction and absorption coefficient for the incident radiation may be deduced. From the absorption integrated over the path of the reflected or transmitted wave (the index of refraction is shown to be large but slowly changing), the emission is inferred by use of Kirchhoff's relation. The principal effect of summing over all γ and integrating over the velocity is to greatly reduce the immense height of the line (when velocity broadening is neglected) without appreciably widening the line. This is intuitively obvious, perhaps, if one considers that the number of electrons having precisely the correct energy and therefore relativistic mass to resonate at a given frequency is reduced as the temperature rises; therefore the height of the line depends inversely on temperature. The effect of electrons resonating below the resonant frequency for a given electron energy, however, cancels the effect of electrons of less energy resonating above the line; therefore the width of the line is essentially unchanged. This "cooperative" cancellation is, we believe, the source of the discrepancy with other work[1,2] which predicts very much greater emission. The calculation was

carried out explicitly for a spatially varying magnetic field characteristic of mirror machines. Even for the fundamental frequency the emission is of the order of 10^{-4} of that for a black body for small laboratory plasmas of 50-kev temperature. The effect is markedly temperature dependent, however, going as $T^{-5/2}$, and rises to 10^{-2} for a 10-kev plasma. Although the spatial fluctuations can be neglected since we are only interested in the coherent transmitted radiation, the fluctuations in velocity space dominate the Doppler motion calculation above 10 kev for electron densities of 10^{14}/cc. We may neglect the fluctuations for higher temperature cases only where the Doppler effect is unimportant, namely for emission perpendicular to the magnetic field and for emission parallel to the field in a mirror machine of less than 20% magnetic field change; for perpendicular motion the fluctuations have a negligible effect below temperatures well in excess of 100 kev. The rth harmonic has an intensity only roughly $r(2kT/mc^2)^{r/2}$ as great as the fundamental, and when detectable furnishes a promising means of determining the plasma temperature by measuring the intensities of the harmonics relative to the fundamental.

*This work was undertaken primarily while the author was a consultant to the University of California Radiation Laboratory at Livermore during 1954-1956.

[1]B. A. Trubnikov and V. S. Kudryavtsev, Proceedings of the Second United Nations International Conference on the Peaceful Uses of Atomic Energy, Geneva, 1958 (United Nations, Geneva, 1958), A/conf. 15/P/2213.

[2]Hayakawa, Hokkyo, Terashima, and Tsuneto, Proceedings of the Second United Nations International Conference on the Peaceful Uses of Atomic Energy, Geneva, 1958 (United Nations, Geneva, 1958), A/conf. 15/P/1330.

[3]J. Schwinger, Phys. Rev. 75, 1912 (1949).

FAST ION HEATING*

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The remarkably fast ion heating observed in high-current "stabilized" gas discharges[1] may be due partly to an ion relaxation process described by Schluter.[2] The periodically varying magnetic field required by this process may arise from the transient excitation of magnetic compression waves[3] occasioned by the violent initial collapse of the discharge. If the axial current flows entirely in a thin surface layer of equilibrium radius r, and the exterior magnetic field is purely azimuthal, the frequencies of the compression modes are given by $\omega_N^2\mu_0\rho r^2 = \sigma^2B^2(1+N^2r^2/\sigma^2R^2)$, where ρ is the average density, B is the average axial (stabilizing) magnetic field, R is the major radius of the torus, and N is zero or an integer. The constant, σ, is a root of $\sigma J_0(\sigma)=J_1(\sigma)$. For compressions that are in phase in a cross section of a discharge, $\sigma=1.841$, the lowest root. In Zeta and Sceptre, $\omega_0/2\pi$ is typically about 0.7 Mc/sec, and the next few modes are only slightly higher in frequency. If we approximate the magnetic field inside a discharge by $B_z=B(1-A\cos\omega t)$, the total energy of the compression mode is of the order $W=\rho\pi^2r^4RA^2\omega^2/8$. In Zeta and Sceptre, under typical conditions, the total energy comes to about $1000A^2$ joules per mode. This is to be compared with the total ion energy, which at 100 ev ion temperature, is roughly 1000 joules in each case.

To get an idea of how quickly the compression mode energy can be given to the ions, we use Schluter's expression for the ion temperature e-folding time, which is $1/\lambda=9(\gamma^2+\omega^2)/A^2\omega^2\gamma$, where γ is the ion collision frequency.[4] We estimate e-folding times on the basis of only one mode, with $A=1$, and for typical operating conditions. In Zeta, for ion temperatures of 1, 10, and 100 ev, we get $1/\lambda$ values of 0.05, 0.005, and 0.1 millisecond. We get corresponding values for Sceptre of 0.3, 0.01, and 0.01 millisecond. In arriving at these figures, we used ion densities of 10^{14} and 9×10^{14} per cm^3, total currents of 1.6 and 1.0×10^5 amp, r values of 14 and 6 cm, and R values of 100 and 57 cm, for Zeta and Sceptre, respectively.

We conclude that this picture is in accord with the observed times and magnitudes of ion heating, with a not unreasonable choice of A values for several of the lower modes of a toroidal discharge.

We wish to acknowledge with thanks the help and support of Dr. H. Motz.

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[1]See the group of papers beginning with P. Thonemann et al., Nature 181, 217 (1958).

[2]A. Schluter, Terzo Congresso Internazionale sui

Fenomeni d' Ionizzazione nei Gas (Società Italiana di Fisica, Milan, 1957), p. 924.

[3]D. Reagan, Phys. Fluids (to be published).

[4]L. Spitzer, Physics of Fully Ionized Gases (Interscience Publishers, Inc., New York, 1956), p. 78.

SPONTANEOUSLY GROWING TRANSVERSE WAVES IN A PLASMA DUE TO AN ANISOTROPIC VELOCITY DISTRIBUTION

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Recently self-excited longitudinal plasma oscillations[1] and self-excited Alfvén waves[2] were found theoretically in plasmas having velocity distributions which deviate from the Gaussian. Upon examining the theories of these phenomena it was found that there exist also self-excited transverse electromagnetic waves, which involve only the electrons of a plasma, provided that their velocity distribution is sufficiently anisotropic. Their existence and rate of growth can be derived from the Boltzmann transport equation, neglecting the collision term and retaining only linear terms of the perturbation:

$$\frac{\partial f}{\partial t} + \vec{v}\cdot\frac{\partial f}{\partial \vec{r}} + \frac{e}{m}\left[\vec{v}\times\vec{B}_0\right]\cdot\frac{\partial f}{\partial \vec{v}} = -\frac{e}{m}\left[\vec{E}+\vec{v}\times\vec{B}\right]\cdot\frac{\partial f_0}{\partial \vec{v}}, \quad (1)$$

where $f_0(\vec{v})$ is a nonisotropic distribution which is stationary in absence of collisions, $\vec{B}_0$ is a constant magnetic field, f is the perturbation of the distribution function, while $\vec{E}$ and $\vec{B}$ represent the perturbation of the electromagnetic field. Assuming that the first order quantities $f(\vec{v},\vec{r},t)$, $\vec{E}(\vec{r},t)$, and $\vec{B}(\vec{r},t)$ depend on $\vec{r}$ and t only through the factor $\exp(i\omega t + i\vec{k}\cdot\vec{r})$, one obtains

$$i(\omega+\vec{k}\cdot\vec{v})f - \frac{e\vec{B}_0}{m}\cdot\left[\vec{v}\times\frac{\partial f}{\partial \vec{v}}\right] = -\frac{e}{m\omega}\left\{\omega\vec{E}\cdot\frac{\partial f_0}{\partial \vec{v}} + \left[\vec{k}\times\vec{E}\right]\cdot\left[\vec{v}\times\frac{\partial f_0}{\partial \vec{v}}\right]\right\}, \quad (2)$$

where the two homogeneous Maxwell equations were used to eliminate $\vec{B}$. It can be seen from (2) that the effects of the anisotropy are contained in the last term of the right-hand side. The magnetic field B_0 is not essential to the growth of disturbances; it is included since it does not complicate the analysis very much. The linear, first order, partial differential equation (2) for f was solved for the following special case:

$$f_0(\vec{v}) = F(v_0, v_z), \quad v_0^2 = v_x^2 + v_y^2,$$

with both $\vec{B}_0$ and $\vec{k}$ parallel to the preferred z-direction. $\vec{E}$ is chosen normal to the vector $\vec{k}$. Using the method of characteristics one obtains after some manipulation

$$f = \frac{e\left\{i(\omega+kv_z)(v_xE_x+v_yE_y) + (eB_0/m)(v_xE_y - v_yE_x)\right\}}{m\omega v_0\left\{(eB_0/m)^2 - (\omega+kv_z)^2\right\}} \times\left\{kv_0\frac{\partial f_0}{\partial v_z} - (\omega+kv_z)\frac{\partial f_0}{\partial v_0}\right\}.$$

Using the two inhomogeneous Maxwell equations together with the relations

$$\delta = e\int f(\vec{v},\vec{r},t)d^3v, \quad \vec{J} = e\int \vec{v}f(\vec{v},\vec{r},t)d^3v, \quad (3)$$

and eliminating the fields $\vec{E}$, $\vec{B}$, one obtains a relation between ω and k:

$$k^2 - \omega^2 = \frac{e^2}{m}\pi\int_{v_0=0}^{\infty}\int_{v_3=-\infty}^{\infty}\frac{(\omega+kv_3)\partial f_0/\partial v_0 - v_0k\,\partial f_0/\partial v_3}{(\omega+kv_3)\pm(eB_0/m)} \times v_0^2\,dv_0\,dv_3. \quad (4)$$

The contour of the v_3-integration must pass above the pole $v_3 = (1/k)(\omega \pm eB_0/m)$ in the v_3-plane.[3] If the distribution function has the form

$$f_0 = \frac{n}{u_0^2u_3(2\pi)^{3/2}}\exp\left[-\frac{v_0^2}{2u_0^2} - \frac{v_3^2}{2u_3^2}\right], \quad (5)$$

it is possible to evaluate the integral and one obtains

$$k^2 - \omega^2 = \omega_p^2\left\{A - \left(A\frac{\omega\pm\omega_c}{u_3k} + \frac{\omega}{u_3k}\right)\phi\left(\frac{\omega\pm\omega_c}{u_3k}\right)\right\}, \quad (6)$$

where

$$\omega_p^2 = ne^2/m, \quad \omega_c = eB_0/m, \quad A = (u_0/u_3)^2 - 1, \quad (7)$$

and

$$\phi(z) = \exp(-\tfrac{1}{2}z^2)\int_{-i\infty}^{z}\exp(\tfrac{1}{2}\xi^2)\,d\xi. \quad (8)$$

Equation (6) for certain real values of k admits of complex values of ω having negative imaginary parts. This can be seen most easily in the case $B_0 = 0$ and in the limit of large ω/u_3k. In this

case $\phi(z) = z^{-1} + z^{-3}$ and (6) becomes

$$\omega^4 - (\omega_p^2 + k^2)\omega^2 - u_0^2 \omega_p^2 k^2 = 0. \qquad (9)$$

The same approximation would have been obtained (in the limit $u_3 k/\omega \ll 1$) for an arbitrary f_0, as may be seen by expanding the denominator of (4) in powers of v_3. The four solutions of (9) are

$$\omega = \pm\{\tfrac{1}{2}[\omega_p^2 + k^2 \pm ((\omega_p^2 + k^2)^2 + 4u_0^2 \omega_p^2 k^2)^{1/2}]\}^{1/2}$$

The solution ω_4, obtained by taking minus signs in both choices, is negative imaginary, showing the existence and rate of growth of self-excited waves. This solution is valid only when $u_3 k/\omega \ll 1$ which implies $u_0 \gg u_3$, as becomes evident from the approximate expression

$$\omega_4 \cong -iu_0 \omega_p k/(\omega_p^2 + k^2)^{1/2}.$$

A more detailed study of (6) shows that for large values of k one obtains damped waves.

The author is indebted to Dr. B. D. Fried for many valuable discussions of these matters.

[1]O. Buneman, Phys. Rev. Lett. 1, 8 (1958).

[2]M. Rosenbluth: "Recent Theoretical Developments in Plasma Stability," paper presented Nov. 1958 at the San Diego Meeting of the Fluid Dynamics Division of the American Physical Society.

[3]N. G. Van Kampen, Physica 23, 641 (1957). This reference contains a comprehensive review of previous literature.

SOLID STATE INFRARED QUANTUM COUNTERS*

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(Received December 29, 1958)

Since maser operation is based on stimulated emission of radiation, masers have an inherent limiting noise temperature of $h\nu/k$ due to spontaneous emission.[1,2] Weber[2] has called attention to the fact that it is possible to construct quantum-mechanical amplifiers without spontaneous emission noise. In fact, this is the usual state of affairs for x-ray or γ-counters. This note describes how a solid counter for infrared or millimeter wave quanta might be constructed in principle.

Consider a crystal containing ions which, among others, have the energy levels shown in Fig. 1. Salts of the rare earths and other transition group ions, which may be embedded as impurities in host lattices, offer examples of this situation.[3,4] The distance between the ground level and level E_2 is such that $h\nu_{12} \gg kT$. If, for example, $h\nu_{12} \sim 100$ cm^{-1} and $T \sim 2°$K, only the ground state is populated. Intense light at the optical frequency ν_{23} is not absorbed, because level E_2 is empty. Whenever an incident infrared quantum $h\nu_{12}$ is absorbed, the light will induce a second transition to E_3, provided its intensity produces transitions at a faster rate than the radiationless decay or spontaneous emission from level E_2 back to the ground state.

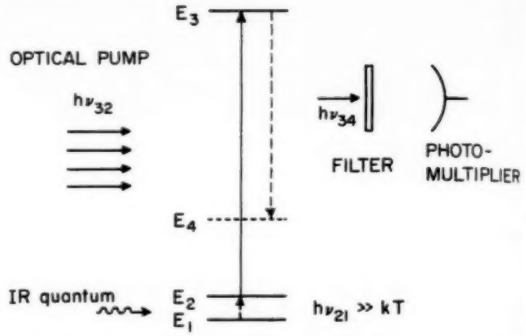

FIG. 1. Infrared quantum counter. Several ions of transition group elements have appropriate energy level diagrams: $h\nu_{21} = 1 - 5000$ cm^{-1}, $h\nu_{32} = 10^4 - 5 \times 10^4$ cm^{-1}.

Spontaneous emission from E_3 to E_2 will produce resonance radiation. The system will be repumped, and several quanta $h\nu_{32}$ may be reemitted for each incident quantum $h\nu_{21}$. It will be difficult to detect these quanta $h\nu_{32}$ in the presence of the intense pumping flux, although one may use discrimination in polarization and direction of propagation. When radiation due to spontaneous emission from level E_3 to E_1 is able to leave the crystal, quanta $h\nu_{31}$ may be counted directly. If this radiation is self-absorbed, a fourth level will provide an effective discrimination in frequency. The fluorescent quanta $h\nu_{34}$ may be counted with a photomultiplier and a suitable filter.

A variation of this scheme is that E_1 is an occupied deep impurity level, E_2 is an empty impurity level, and E_3 represents the conduction band. The incident quantum $h\nu_{21}$ triggers a photoconductive avalanche in the semiconductor near absolute zero of temperature.

It is illuminating to point out the relationship with optical pumping methods proposed by

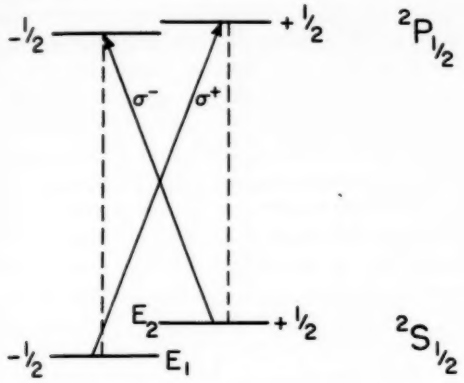

FIG. 2. Optical pumping between $^2S_{1/2}$ and $^2P_{1/2}$ doublets (reference 5). σ^+ pumping produces a maser at $h\nu_{21}$, the ground-state splitting. σ^- pumping produces a quantum counter at $h\nu_{21}$.

Kastler[5] and with three-level masers.[6,7] Consider the σ^- and σ^+ transitions between the $^2S_{1/2}$ and $^2P_{1/2}$ states shown in Fig. 2. It makes no difference whether these transitions are separated by the sense of circular polarization[5] or by the Zeeman frequency splitting.[8] If the σ^+ transition is pumped, one obtains maser action in the ground-state doublet. If the σ^- transition is pumped, one prepares a quantum counter. To eliminate thermal noise in this counter and keep all particles in the ground state E_1, either free atoms in a beam or a crystal near absolute zero of temperature should be used.

The condition on the intensity for optical pumping is the same to achieve maser action or quantum counting. Using the parlance of magnetic resonance, one has to approach saturation of the optical absorption coefficient at the pumping frequency. The optically induced transitions must have a faster rate than the competing relaxation mechanisms between levels E_2 and E_1. In favorable cases it appears possible to produce the required optical intensities.

For the quantum counter, details of the decay from level E_3 are not important, provided radiationless transitions are not dominant. For the maser action, one requires a decay to level E_2 at the faster rate than the relaxation rate from E_2 to E_1. It is evident that in either case more than one optical transition may be pumped.

Although the prospect of reaching a quantum efficiency of 100% or even 10% appears remote, the type of infrared detector proposed here may be useful even if it falls short by many orders of magnitude of its ultimate potential.

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*First communicated at the Symposium on New Techniques in Microwave and Radio-frequency Physics, National Academy of Sciences, Washington, D. C., April, 1958 (unpublished).

[1]Gordon, Zeiger, and Townes, Phys. Rev. 99, 1264 (1955); R. V. Pound, Ann. Phys. 1, 24 (1957); M. W. Muller, Phys. Rev. 106, 8 (1957); M. W. P. Strandberg, Phys. Rev. 106, 617 (1957); Shimoda, Takahasi, and Townes, J. Phys. Soc. Japan 12, 686 (1957).

[2]J. Weber, Phys. Rev. 108, 537 (1957).

[3]G. H. Dieke and L. Heroux, Phys. Rev. 103, 1227 (1956); G. H. Dieke and L. Leopold, J. Opt. Soc. Am. 47, 944 (1957).

[4]S. Sugano and Y. Tanabe, J. Phys. Soc. Japan 13, 880 (1958); S. Sugano and I. Tsujikawa, J. Phys. Soc. Japan 13, 899 (1958).

[5]A. Kastler, Proc. Phys. Soc. London A67, 853 (1954); J. Opt. Soc. Am. 47, 460 (1957).

[6]N. G. Basov and A. M. Prokhorov, J. Exptl. Theoret. Phys. U.S.S.R. 27, 431 (1954).

[7]N. Bloembergen, Phys. Rev. 104, 324 (1956).

[8]H. Bucka, Z. Physik 141, 49 (955).

EXTENDED FINE STRUCTURE OF *K* X-RAY ABSORPTION EDGES IN PERMINVAR

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In other communications,[1-3] the structural origin and mechanism of the magnetic annealing of soft, magnetic alloys have been investigated. It has been found that the ability of soft, magnetic alloys containing Fe, Co, and Ni to exhibit magnetic annealing properties is due to oxygen present as an impurity. Single crystals with an oxygen content of 0.001-0.002% respond to field heat treatment. Single crystals with an oxygen content of 0.0001% fail to respond to field heat treatment.

It is well known that x-ray spectra, both in emission and absorption, are sensitive to the state of chemical combination of the elements under study. Thus there is a possibility that a study of the x-ray spectra of Fe, Co, and Ni in soft, magnetic alloys might yield information on how the metal atoms are associated with the oxygen impurity.

The extended fine structure of the *K* x-ray ab-

sorption edges of Fe, Co, and Ni has been investigated in polycrystalline Perminvar foils (30 Fe, 25 Co, 45 Ni). Torque magnetometer data[3] for the Perminvar alloys show a response to field heat treatment that is several times greater than that exhibited by other soft, magnetic alloys such as the Permalloys. The oxygen content in Perminvar foils is considerably higher than in single crystals, about 0.02%.

The absorption spectra have been obtained with a General Electric diffractometer operated as a tube spectrometer. A tungsten target x-ray tube with a line focus of 0.8×15 mm was used at a 5° take-off angle giving a projected line focus width of about 0.07 mm. A 0.4° beam slit, a LiF analyzing crystal, and a 0.02° detector slit completed the x-ray optics. The detector was a gas flow proportional counter with a 90% argon, 10% methane gas mixture. The absorbing specimens were metal tapes which were thinned down to about 0.00025 in.

In Fig. 1 the extended fine structure of the K absorption edge of Ni in a 30-25-45 Perminvar and in Ni metal is illustrated. The ordinates are expressed as $\ln(I_0/I)$, where I_0 is the incident beam intensity and I is the transmitted beam intensity. They are therefore proportional to μt, where μ is the linear absorption coefficient and t is the absorber thickness. The abscissas are expressed in electron volts with a zero of energy taken at the point where the value of μ has attained half its initial increase from the low-energy side of the absorption edge to the top of the first maximum. The extended fine structure for Ni metal is well known and is observed for all elements that crystallize in cubic or hexagonal closest packing. The same extended fine structure is observed for elements in solid solution in closest packed alloys. Indeed we have observed this same extended fine structure for Fe in 30-25-45 Perminvar, 20-60-20 Perminvar, Permalloy, and 18-8 stainless steel; and for Co in Co metal, 30-25-45 Perminvar, and 20-60-20 Perminvar. However, the extended fine struc-

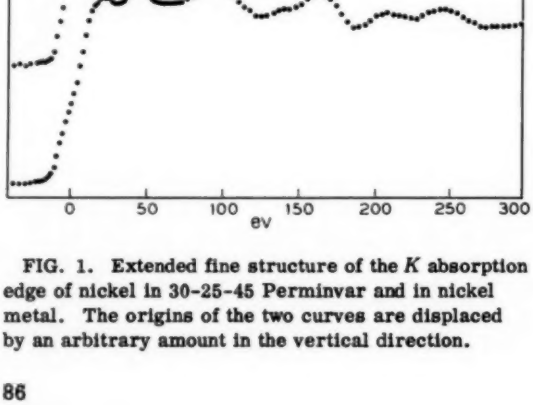

FIG. 1. Extended fine structure of the K absorption edge of nickel in 30-25-45 Perminvar and in nickel metal. The origins of the two curves are displaced by an arbitrary amount in the vertical direction.

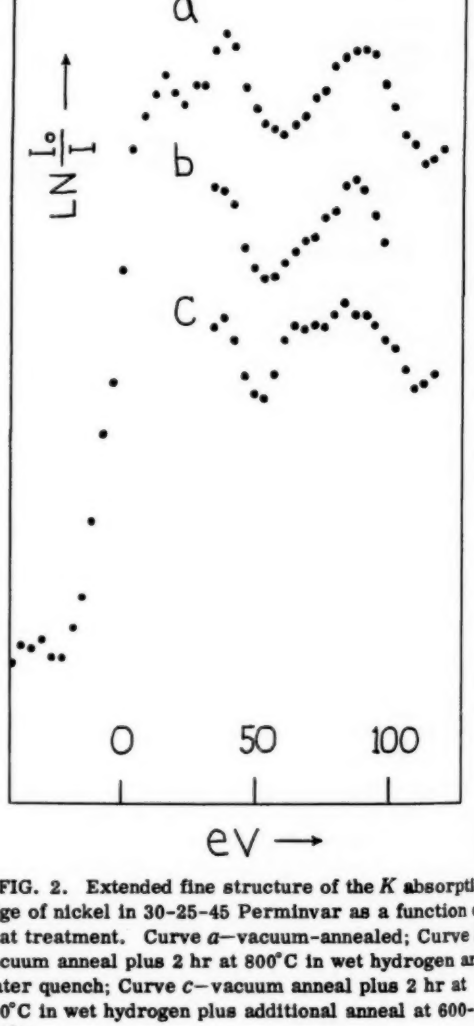

FIG. 2. Extended fine structure of the K absorption edge of nickel in 30-25-45 Perminvar as a function of heat treatment. Curve a—vacuum-annealed; Curve b—vacuum anneal plus 2 hr at 800°C in wet hydrogen and water quench; Curve c—vacuum anneal plus 2 hr at 800°C in wet hydrogen plus additional anneal at 600-400°C.

ture for Ni in 30-25-45 Perminvar which has been suitably annealed has a pronounced difference. At about 65 ev above the absorption edge there is an extra maximum which is comparable in magnitude with any of the principal maxima in the structure. No indication of the new, anomalous peak has ever been observed in the Fe and Co edges of Perminvar nor has it yet been seen in the Ni edge of the Permalloys. The occurrence of the anomalous absorption peak at about 65 ev is directly connected with a particular state of aggregation of oxygen in the alloy. Curves *a*, *b*, and *c* of Fig. 2 show a region of the extended fine structure where the anomalous peak is deliberately and controllably introduced through heat treatment in wet hydrogen. Curve *a* is for vacuum-annealed Perminvar and is typical of f.c.c. nickel. Curve *b* results from 2 hr at 800°C in wet hydrogen with a water quench. Curve *c* shows the anomalous peak after additional annealing at 600-400°C. The temperature range where the new peak is strongly developed is identical with the temperature range for magnetic annealing.

Transmission electron diffraction results for the three states illustrated in Fig. 2 all exhibit a precipitated metal monoxide. It must be concluded that precipitated oxides are not the cause of the anomalous peak but rather that an earlier state of segregation is responsible. The present indications are that the "oxygen faults" described in reference 2 are the cause of the anomalous absorption in the Ni edge as well as the magnetic annealing properties.

The authors are indebted to E. A. Nesbitt for aid in the preparation of alloys and their heat treatments.

[1]R. D. Heidenreich and E. A. Nesbitt, Phys. Rev. 105, 1678 (1957).
[2]Heidenreich, Nesbitt, and Burbank, J. Appl. Phys. (to be published).
[3]E. A. Nesbitt and R. D. Heidenreich, J. Appl. Phys. (to be published).

INVERSION SYMMETRY OF THE *M* CENTER

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(Received January 9, 1959)

According to the model proposed by Seitz,[1] *M* centers in alkali halide crystals consist of two negative-ion vacancies, one positive-ion vacancy, and an electron, as shown schematically in Fig. 1 (a). The low symmetry (C_{2v}) of this model has made the *M* center a particularly interesting subject for theoretical[2] and experimental[3-6] study. The purpose of the present note is to suggest a slight modification of the model which has all the essentials of the original but which, in addition, exhibits inversion symmetry.

Consider the ionic configuration of Fig. 1 (a) with the electron removed. If the three vacancies remain associated, the positive ion marked *M* will trade places with the positive-ion vacancy at a certain jump frequency determined in part by the nature of the potential barrier separating these two equivalent positions. When *M* is in the region midway between these two positions, it experiences practically no repulsive force from the ions in the plane of the center, and the elec-

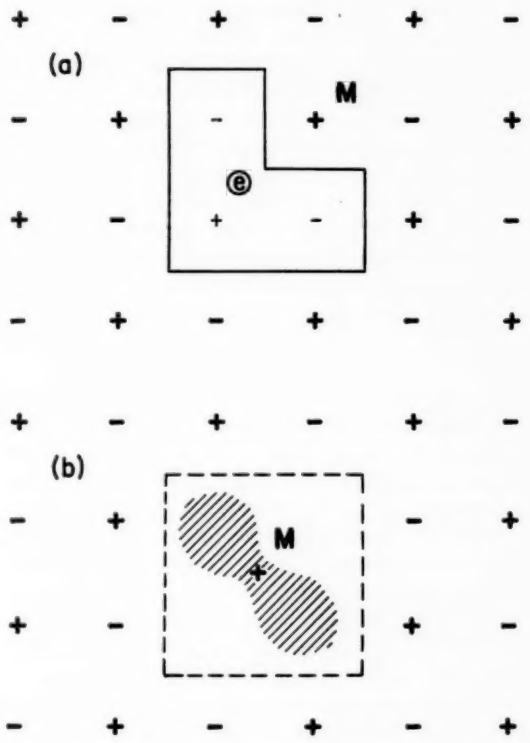

FIG. 1. Models of the *M* center in alkali halide crystals. (a) The original Seitz model, (b) a proposed revision of its configuration. In (a) the boxes represent vacancies and *e* represents an electron. In (b) the region outlined by a dashed line is occupied only by the ion *M* and the electron, whose possible distribution is indicated by the shaded region.

trostatic potential has a saddle point at the midpoint. With the electron present, it is reasonable to assume that the jumping still takes place but that the electron, in trying to adjust adiabatically to the configuration of the center, associates itself as closely as possible with the positive ion. The effect of the electrostatic saddle point is thus at least partially washed out and the ion may spend more of its time in the central position, as sketched in Fig. 1 (b). Once this configuration is established the electron will be drawn toward the positively charged corners of the "pillbox" in which it now finds itself. Thus the electronic, as well as configurational, symmetry of the center is changed from C_{2v} to D_{2h}. One may consider the central ion-electron pair as simply an alkali atom, distorted toward the nearest neighboring positive ions.

Overhauser and Rüchhardt[5] tentatively conclude that centers responsible for the M absorption band have inversion symmetry on the basis of the absence of a large differential Stark effect as measured by line broadening at high external electric fields. Centers whose configurations are that of Fig. 1 (a) have permanent dipole moments and such a Stark effect is likely, according to these authors' computations. Jacobs[6] was unable to find dielectric loss which could be associated with jumping of the M-center positive ion across a sizeable potential barrier. The model of Fig. 1 (b) is suggested in an attempt to reconcile these negative results with the original model, which is a plausible result of coagulation of F centers and neutral vacancy pairs.[1] Since the basic composition and the predominantly (110) symmetry of the center have been left unchanged, no reappraisals of experiments on formation and optical bleaching of the center[1-3] are necessary, while the lack of a sizable potential barrier and the presence of inversion symmetry seem to be complementary to each other on the revised model.

Detailed computations are necessary to establish the quantitative plausibility of the inversion symmetry model, but spin resonance experiments will probably be of more immediate assistance in its evaluation. A resonance associated with M centers has already been identified,[4] but this resonance is hardly resolved from that of F centers and double resonance measurements[7] (presently in progress at the University of Illinois[8]) are necessary to obtain more definite information on the center's structure.

The author is indebted to Dr. G. Jacobs for interesting discussions of his experiments and to Professor F. Seitz for his kind interest in this work.

[1]F. Seitz, Revs. Modern Phys. 26, 7 (1954).
[2]Inui, Uemura, and Toyozawa, Progr. Theoret. Phys. (Japan) 8, 355 (1952); B. S. Gourary and P. J. Luke, Phys. Rev. 107, 960 (1957).
[3]M. Ueta, J. Phys. Soc. Japan 7, 107 (1952); J. Lambe and W. D. Compton, Phys. Rev. 106, 684 (1957); W. D. Compton and C. C. Klick, Phys. Rev. 112, 1620 (1958). H. Kanzaki, Phys. Rev. 110, 1063 (1958).
[4]N. W. Lord, Phys. Rev. 106, 1100 (1957); H. Kawamura and K. Ishiwatari, J. Phys. Soc. Japan 13, 33 (1958).
[5]A. W. Overhauser and H. Rüchhardt, Phys. Rev. 112, 722 (1958).
[6]G. Jacobs, J. Chem. Phys. 27, 218 (1957).
[7]G. Feher, Phys. Rev. 105, 1122 (1957).
[8]C. P. Slichter (private communication).

SPIN-LATTICE RELAXATION FROM STATE OF NEGATIVE SUSCEPTIBILITY*

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A state of negative magnetic susceptibility has been demonstrated in potassium chromicyanide by using a 180° pulse technique. The spin-lattice relaxation from this state has been observed in the time domain, and the complex susceptibility shows no change in slope as it passes through zero.[1]

Negative paramagnetic susceptibility has been demonstrated in nuclear magnetic resonance with the use of both rapid passage[2] and 180° pulse techniques.[3] Negative electron paramagnetic susceptibility has been produced in the three-level maser[4] and in the two-level maser by means of rapid passage.[5,6] The feasibility of generating a negative susceptibility in semiconductors by using spin-echo techniques has also been demonstrated.[7] We have produced negative electron magnetic resonance susceptibility by means of a 180° pulse. The working material was $K_3Co(CN)_6$ containing 0.1% Cr^{3+}. The $-\frac{1}{2}\rightarrow+\frac{1}{2}$ transition at 9000 Mc/sec in a field of 3150 gauss, oriented parallel to the crystalline c-axis, was used.

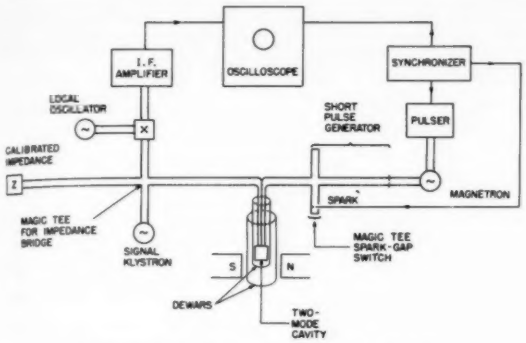

FIG. 1. Block diagram of equipment.

The equipment has been described elsewhere.[8] The general arrangement is shown schematically in Fig. 1. A microwave pulse, roughly 10 mμsec in duration, is generated from energy stored in a cavity formed from a length of waveguide excited by an 0.5-μsec magnetron pulse. One end of the waveguide is opened by firing a spark which unbalances a magic tee "bridge" by changing a plane of reflection. The sample is placed in a cavity that possesses two degenerate, orthogonal modes. The mode to which the short pulse is applied has a very low Q (roughly 30). The other mode, which has a high Q, is used to observe the magnetic resonance of the sample. In order to use a small measuring signal, a superheterodyne detection system (30-Mc/sec i.f.) is used. An impedance bridge is used to measure the instantaneous complex impedance of the cavity.

The quantity plotted on the ordinate in Fig. 2 is proportional to the magnetic susceptibility. The cavity filling factor was not determined. The point at which the susceptibility passes from negative to positive values is indicated. It is interesting to note that there is no change in slope at this point. The observed T_1 is 8.65 msec.

As a check on the tuning of the cavity and magnetic field and on the correct interpretation of the experiment, it was also possible to display the complex reflection coefficient of the cavity on an oscilloscope by means of auxiliary equipment that is not shown in Fig. 1 (i.e., a high-speed Smith-chart plotter).

To our understanding, a simple relaxation process uncomplicated by anomalous phonon effects is to be expected in crystals that are as dilute as this at 4.2°K. This result indicates that spin spectral diffusion[9] and the effect of the other levels do not strongly influence the simple decay of the susceptibility. That anomalies can occur in less dilute crystals can be seen, since others,[10] using saturation techniques, obtain for this crystal with 0.5% Cr^{3+} a T_1 of 0.2 sec, independent of spin quantum numbers. Presumably, these measurements have determined a lattice-bath relaxation time rather than a true spin-lattice

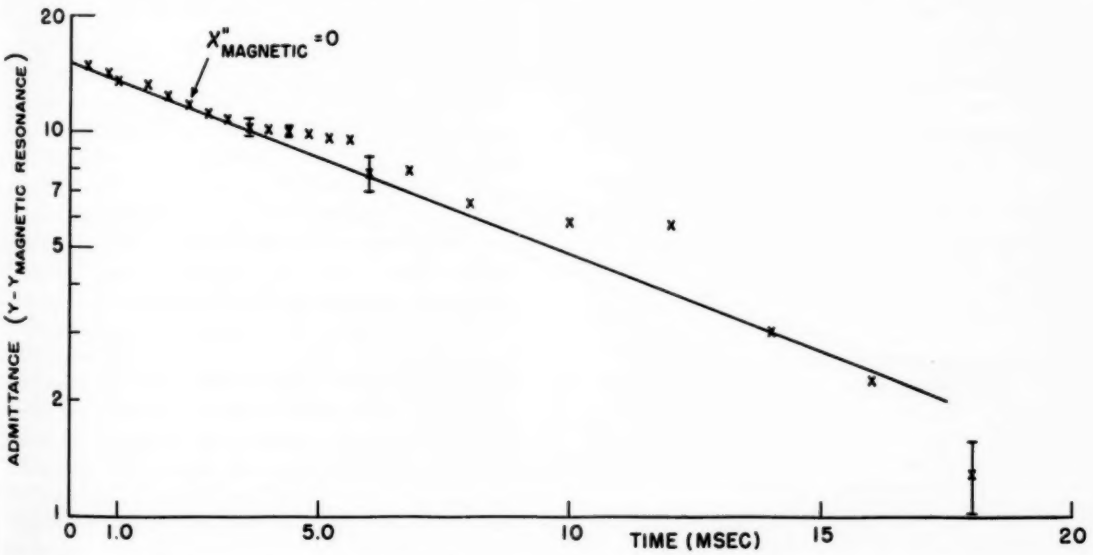

FIG. 2. Semilogarithmic plot of a quantity proportional to magnetic susceptibility versus time.

relaxation time.[11]

Much remains to be done experimentally to develop a sound understanding of the spin-lattice relaxation process, but we feel that this demonstration of negative susceptibility by 180° pulse techniques, and the direct observation of the associated relaxation time from a state of negative magnetic susceptibility shed new light on spin-lattice relaxation processes in paramagnetic crystals.[12]

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†Now with Sperry Gyroscope Company, Great Neck, New York.

[1]Research on Paramagnetic Resonances, Third Quarterly Progress Report, Signal Corps Contract DA36-039-sc-74895, January - April 1958. (The complete research report, of which this work is only a part, has been unduly delayed, and will be presented later.)

[2]Bloembergen, Purcell, and Pound, Phys. Rev. 73, 679 (1948).

[3]E. L. Hahn, Phys. Rev. 80, 580 (1950).

[4]Scovil, Feher, and Seidel, Phys. Rev. 105, 762 (1957); Strandberg, Davis, Faughnan, Kyhl, and Wolga, Phys. Rev. 105, 447 (1957).

[5]Feher, Gordon, Buehler, Gere, and Thurmond, Phys. Rev. 109, 221 (1958).

[6]Chester, Wagner, and Castle, Phys. Rev. 110, 281 (1958).

[7]J. P. Gordon and K. D. Bowers, Phys. Rev. Lett. 1, 368 (1958).

[8]Research on Paramagnetic Resonances, First Quarterly Progress Report, Signal Corps Contract DA36-039-sc-74895, July - October 1957 (unpublished).

[9]A. M. Portis, Phys. Rev. 104, 584 (1956).

[10]A. L. McWhorter and J. W. Meyer, Phys. Rev. 109, 312 (1958).

[11]M. W. P. Strandberg, Phys. Rev. 110, 65 (1958).

[12]Similar results were reported by Castle, Chester, and Wagner at the Molecular Electronics Symposium held at the Statler-Hilton Hotel, Washington, D. C., November 13-14, 1958.

COLLECTIVE EXCITATIONS AND THE MEISSNER EFFECT*

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The theory of superconductivity of Bardeen, Cooper, and Schrieffer[1] has been generalized by Bogoliubov, Tolmachov, and Shirkov[2] and by Anderson[3] to include collective excitations of the electrons. Both make use of the random phase approximation. We have calculated the response of a superconductor at $T=0°K$ to a static magnetic field described in a general gauge, by use of Anderson's method. As he points out, the theory satisfies the sum rules and therefore should be gauge invariant. Aside from a small correction from the transverse collective modes, results are the same as those of BCS, who used the gauge, $\mathrm{div}A=0$, and included only quasi-particle excitations.

Longitudinal collective excitations (plasma oscillations if Coulomb interactions are included) do not play a direct role if $\mathrm{div}A=0$ but must be considered if the calculation is performed in a general gauge. Their crucial part in the satisfaction of the longitudinal sum rule has been discussed by Anderson.[4] Pines and Schrieffer[5] have shown that the plasmons shield the electrons in such a way that quasi-particles do not contribute to the longitudinal current so that the latter is described entirely in terms of the plasmon variables.

In the present calculation it is assumed that the normal-state wave functions are plane waves of energy $\epsilon_k=(\hbar^2k^2/2m)-\epsilon_F$ and that the superconducting transition is caused by a two-particle interaction

$$\sum_{k,k'} V(k,k')\,C_{k'\sigma'}{}^{*}\,C_{-k'+q,\sigma}{}^{*}\,C_{-k+q,\sigma}C_{k\sigma'}.$$

To ensure the gauge invariance of the results, $V(k,k')$ is a function of $\vec{k}-\vec{k}'$, only.

One should really start with the original gauge-invariant Hamiltonian in which the electron-phonon interaction is still present. Wentzel[6] has attempted to calculate the Meissner effect from Fröhlich's Hamiltonian[7] in a completely gauge-invariant manner but his procedure does not introduce the collective excitations. His result disagrees with ours in the London limit. To get the correct result by his method it would presumably be necessary to sum, to all orders in the interaction V, appropriate terms in the perturbation expansion. This question is discussed by Pines and Schrieffer.[8] The author[9] has calculated the effect from the same Hamiltonian in the gauge $\mathrm{div}A=0$ and has obtained the same result as BCS. This calculation shows that Fröhlich's Hamiltonian leads to essentially the same results as that involving the two-particle interaction.

For our interaction the energy gap parameter is I_k, the solution of the integral equation

$$I_K = -\tfrac{1}{2}\sum_k V(K,k)\, I_k/E_k, \qquad E_k = (\epsilon_k^2 + I_k^2)^{1/2}.$$

Near the Fermi surface I_k is substantially constant and corresponds to ϵ_0 of BCS. In the presence of a static vector potential $\vec{a}(q)\exp(i\vec{q}\cdot\vec{r})$ the current density is

$$\vec{j}(q) = \frac{e\hbar}{2m(2\pi)^3}\sum_{k,\sigma} C_{k+q,\sigma}^{*} C_{k,\sigma}(2\vec{k}+\vec{q}) - \frac{e^2}{mc\Omega}\sum_{k,\sigma} C_{k+q,\sigma}^{*} C_{k,\sigma}.$$

From the equations of motion of products such as $C_{k,\sigma}^{*} C_{k+q,\sigma}$ and $C_{k\uparrow} C_{-k\downarrow}$ we find for the expectation value of this current density in the ground state

$$\vec{j}(q) = \frac{e^2\hbar^2}{4m^2c\Omega}\left\{\sum_k (2\vec{k}+\vec{q})\,\frac{\vec{a}\cdot(2\vec{k}+\vec{q})}{E_k+E_{k+q}}\left(1-\frac{\epsilon_k\epsilon_{k+q}+I_kI_{k+q}}{E_kE_{k+q}}\right) - \tfrac{1}{2}\sum_k (2\vec{k}+\vec{q})\,A_k^{\,q}\,\frac{\epsilon_k I_{k+q} - \epsilon_{k+q} I_k}{E_kE_{k+q}(E_k+E_{k+q})}\right\} - \frac{ne^2}{mc}\vec{a}(q),$$

where $A_k^{\,q}$ is the expectation value of a collective coordinate and is given by the integral equation

$$A_K^{\,q} + \sum_k V(K,k)\,\tfrac{1}{2}\left(1+\frac{\epsilon_k\epsilon_{k+q}+I_kI_{k+q}}{E_kE_{k+q}}\right)\frac{A_k^{\,q}}{E_k+E_{k+q}} = \sum_k V(K,k)\,\frac{(\epsilon_k I_{k+q} - \epsilon_{k+q} I_k)}{E_kE_{k+q}(E_k+E_{k+q})}\,\vec{a}\cdot(2\vec{k}+\vec{q}).$$

This result is formally independent of whether or not we include the Coulomb interactions. For a field described by a static longitudinal vector potential, we can write $\vec{a} = \vec{q}\phi$ and we find $A_k^{\,q} = (4m/\hbar^2)(I_k + I_{k+q})\phi$. The final integration for $j(q)$ can be carried out exactly and leads to $j(q) = 0$, as it should in a gauge-invariant theory.

For a transverse field, we find in the London limit, $q \to 0$, that the integrands tend to zero and, as there is no singularity, $j(0) = -(ne^2/mc)a(0)$. We see here quite clearly the different parts played by the collective coordinates in longitudinal and transverse fields. In the Pippard limit, $q\xi_0 \gg 1$, there is a correction to the BCS result arising from the collective coordinates. This correction depends strongly on the dependence of $V(K,k)$ on θ, the angle between k and K, but not on its energy dependence. If $V(K,k)$ were independent of θ, the correction would be zero. We have estimated the magnitude of the correction by choosing

$$V(k,K) = -\tfrac{3}{4}\overline{V}(1-\cos\theta)^2.$$

In this case

$$j_\perp(q) = \frac{-3c}{16q\xi_0\lambda_L^{\,2}(0)}\left\{1 - \frac{16}{\pi^3 q\xi_0}\ln(\pi q\xi_0) - \frac{8N(0)\overline{V}}{3\pi^3 q\xi_0}\left[\ln(\pi q\xi_0) - \tfrac{4}{3}\right]^2\right\} a_\perp.$$

The first two terms are those given by BCS while the third is the new correction. For $\pi q\xi_0 = 10$ and $N(0)\overline{V} = 0.3$, the ratio of the third term to the second is 0.02, which suggests that the correction is small. As the correction is sensitive to the dependence of the potential on angle, the argument is not conclusive.

Fuller details of this calculation together with a calculation of the dielectric constant of a superconductor within the random phase approximation will be published shortly. It is a pleasure to acknowledge the advice of Professor Bardeen at all stages of this work.

*This work was supported in part by the Office of Ordnance Research, U. S. Army.

[1]Bardeen, Cooper, and Schrieffer, Phys. Rev. 108, 1175 (1957), hereafter referred to as BCS.

[2]Bogoliubov, Tolmachov, and Shirkov, Joint Institute for Nuclear Research, Dubna, 1958 (to be published).

[3]P. W. Anderson, Phys. Rev. (to be published).

[4]P. W. Anderson, Phys. Rev. 110, 827 (1958).

[5]D. Pines and J. R. Schrieffer (to be published).

[6]G. Wentzel, Phys. Rev. 111, 1488 (1958).

[7]H. Fröhlich, Proc. Roy. Soc. (London) A215, 291 (1952).

[8]D. Pines and J. R. Schrieffer, Phys. Rev. Lett. 1, 407 (1958).

[9]G. Rickayzen, Phys. Rev. 111, 817 (1958).

INFLUENCE OF COLD WORK ON THE RESISTIVITY OF DILUTE COPPER ALLOYS*

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(Received January 5, 1959)

The influence of plastic deformation and subsequent annealing on the electrical properties of pure copper has been studied extensively in recent years.[1] Although work on several alloys has been reported,[2] it appears that no systematic

investigations of an alloy series have been undertaken. We have begun a study of copper, silver, and gold alloyed in dilute concentrations with their neighbors in the periodic table. We report here some preliminary results which, we feel, are sufficiently unusual to warrant presentation prior to the completion of the entire program.

The following alloys of copper—CuGe, CuIn, CuSn, and CuSb—containing one atomic percent solute were prepared by melting the constituent metals in vacuum. The copper used in the preparation was 99.999% pure; the purity of the solute metals was 99.99% or better. The ingots, cylinders of approximately 5 mm diameter, were then severely deformed at room temperature by cold-rolling into strips 5 mils thick and about 2 cm wide. From these strips we cut narrow ribbons, roughly 2 mm wide, which were used in resistance and thermoelectric power measurements. Throughout the preparation of the samples they were carefully cleaned by etching in hydrochloric and nitric acids.

An indication of the macroscopic homogeneity of the samples was obtained from thermoelectric power data. The thermoelectric emf was found to be independent of the region along the sample within which the temperature gradient was maintained, being a function only of the temperature at the junctions.[3]

The resistance ratio $p = R_{4.2°}/(R_{273°} - R_{4.2°})$ of the samples before and after a variety of different annealing procedures is shown in Table I. We have also listed the residual resistivities $\Delta\rho$, calculated under the assumption that the ideal resistivity of pure copper is 1.55 μohm-cm at 273°K.[4]

The resistance ratios of the annealed copper samples show that the copper was of high purity. The difference in the residual resistivities of the annealed and unannealed pure copper is of the expected magnitude.[1] The dislocation density in the cold-worked pure copper is probably in the neighborhood of 10^{11} lines/cm^2.[5]

The striking feature of our results is the increase of p and of $\Delta\rho$ of each of the alloys as a consequence of annealing. In those cases in which the samples were annealed at a lower temperature prior to the high-temperature anneal, a sizeable increase in the residual resistivity took place during the first anneal. Also, annealing at a lower temperature following the high-

Table I. Resistance ratio $p = R_{4.2°}/(R_{273°} - R_{4.2°})$ and residual resistivity $\Delta\rho$ of cold-worked and annealed copper and copper alloys.

Sample	Annealing treatment	p	$\Delta\rho$ μohm-cm
Cu-I	none	5.4×10^{-3}	8.4×10^{-3}
Cu-I	3 hr, vacuum, 950°C	1.85×10^{-3}	2.9×10^{-3}
Cu-II	none	5.1×10^{-3}	7.9×10^{-3}
Cu-II	3 hr, 1 atmos He, 850°C	2.7×10^{-3}	4.2×10^{-3}
CuGe-I	none	1.57	2.43
CuGe-I	8 hr, vacuum, 500°C	1.97	3.1
CuGe-I	additional 8 hr, vacuum, 600°C	2.01	3.2
CuGe-I	additional 3 hr, vacuum, 950°C	2.03	3.2
CuIn-I	none	0.387	0.60
CuIn-I	3 hr, 1 atmos He, 950°C	0.57	0.88
CuSn-I	none	1.1	1.7
CuSn-I	8 hr, vacuum, 450°C	1.29	2.0
CuSn-II	none	1.1	1.7
CuSn-II	3 hr, 1 atmos He, 950°C	1.735	2.68
CuSn-II	additional 8 hr, 1 atmos He, 450°C	1.543	2.4
CuSb-I	none	0.85	1.32
CuSb-I	8 hr, vacuum, 450°C	2.59	4.0
CuSb-II	none	0.86	1.33
CuSb-II	3 hr, 1 atmos He, 950°C	2.82	4.37
CuSb-II	additional 8 hr, 1 atmos He, 450°C	2.78	4.31

temperature anneal did not significantly alter the residual resistivity. The only samples which behaved in the expected manner were the pure copper samples; these did show a decrease in residual resistivity as a result of annealing after cold work.

Our results cannot be attributed to (1) diffusion of silicon or other impurities from the Vycor tube, in which the samples were contained during the anneal, into the samples; (2) diffusion of impurities present in traces in the helium gas into the samples during anneal; (3) quenching of vacancies into the samples.

(1) If such contamination had occurred, the resistivities of the pure copper samples should also have increased upon annealing. Moreover, annealing at approximately 500°C should have been much less effective than annealing at about 900°C.

(2) Prior to admitting the helium gas into the tube housing the sample it was passed through a cold trap containing activated charcoal. This procedure should have trapped most chemically active impurities that might have been present in the gas. However, more convincing than such arguments is the fact that the resistivity changes are the same for samples annealed in vacuum and in helium. Consequently, the effects which have been observed cannot be attributed to the presence, or absence of helium.

(3) The samples were quenched by dipping the Vycor tube into water. The quenching rate, and the final temperature, were such that no significant residual vacancy concentration could remain.

The explanation of our result which we suggest is the following. During the course of cold work dislocations are produced, and their motion generates a large number of vacancies,[6] far in excess of the thermal equilibrium concentration of these defects. Since vacancies in copper are mobile at room temperature, these excess vacancies provide a mechanism for diffusion of solvent and solute atoms at a rate far greater than that normally observed at room temperature. For the alloys studied, the solutes would tend to diffuse more readily than the solvent atoms.[7] We suggest that the impurity atoms are trapped at dislocation lines, which in copper are split into two partial dislocations each, and at the stacking faults between these partial dislocations.[8] A dislocation density near 10^{12} lines/cm^2 would provide enough sites to trap all of the impurities at the dislocations and stacking faults. This density is somewhat higher than is usually postulated for plastically deformed pure copper. However, it is probable that in alloys, where impurity atoms can pin dislocations, the dislocation density in severely deformed samples is higher than in pure copper.

When the samples are annealed at high temperature the impurities are freed from their traps and, furthermore, the number of dislocations is greatly diminished. Following the anneal, therefore, the distribution of impurities is microscopically homogeneous in contrast to the microscopically inhomogeneous distribution which prevailed after cold work.

In the well-annealed samples the impurities scatter as independent point defects. In the cold-worked samples the scattering of conduction electrons is reduced for one or both of the following reasons:

(a) Lines and planes of impurities, randomly oriented, lead to a reduction of scattering by a factor of $\frac{1}{3}$ to $\frac{2}{3}$ when one averages over orientations of the lines and planes.

(b) Impurities whose valence exceeds that of the solvent atoms present a scattering cross section which is diminished by local lattice expansion.[9] These same impurities, which normally expand the lattice, will naturally locate themselves at the portion of the dislocation line opposite the extra half-plane, where the lattice is already expanded. Here the local effective lattice expansion is significantly greater than it is in the neighborhood of the same impurity when it is in a substitutional position in the lattice. Consequently, the scattering cross section of an impurity trapped at a dislocation line is less than it is for the same impurity at a substitutional site of the lattice.

The results reported here are not the first which show that in some cases plastic deformation of alloys leads to a reduction in residual resistivity.[2] Of particular interest in this connection are the results of Weyerer[10] who observed a decrease in resistivity of commercial copper after deformation. Further studies of the resistivity of dilute noble metal alloys following plastic deformation and subsequent anneal are now in progress.

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*Supported in part by the Office of Ordnance Research, U. S. Army.

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[1]C. W. Berghout, Acta Metallurgica 4, 212 (1956);

H. G. van Bueren, Acta Metallurgica 3, 524 (1955); D. J. van Ooijen, dissertation, University of Delft, 1956 (unpublished); Blewitt, Coltman, and Redman, Report of a Conference on Defects in Crystalline Solids (The Physical Society, London, 1955), p. 369; J. Molenaar and W. H. Aarts, Nature 166, 690 (1950); M. J. Druyvesteyn and J. A. Manintveld, Nature 168, 868 (1951); W. H. Aarts and R. H. Jarvis, Acta Metallurgica 2, 87 (1954); R. H. Pry and R. W. Henning, Acta Metallurgica 2, 318 (1954); T. Broom, Proc. Phys. Soc. (London) B65, 871 (1952).

[2]Clarebrough, Hargreaves, and West, Proc. Roy. Soc. (London) A232, 252 (1955); J. O. Linds, Appl. Sci. Research B4, 73 (1956).

[3]R. H. Kropschot and F. J. Blatt (to be published).

[4]E. Grüneisen, Ergeb. exakt. Naturw. 21, 50 (1945).

[5]W. A. Harrison, J. Phys. Chem. Solids 5, 44 (1958).

[6]F. Seitz, Advances in Physics, edited by N. F. Mott (Taylor and Francis, Ltd., London, 1953), Vol. 1, p. 43.

[7]C. T. Tomizuka and L. Slifkin, Phys. Rev. 96, 610 (1954).

[8]A. H. Cottrell, Dislocations and Plastic Flow in Crystals (Oxford University Press, Oxford, 1953); H. Suzuki, Sci. Repts. Research Insts., Tohoku Univ. Ser. A 4, 455 (1952); H. Suzuki, Dislocations and Mechanical Properties of Crystals (John Wiley and Sons, New York, 1957), p. 361.

[9]F. J. Blatt, Phys. Rev. 108, 285 (1957).

[10]H. Weyerer, Z. Metallk. 44, 51 (1953).

OBSERVATIONS OF HYDROGEN VIBRATION FREQUENCIES IN PHOSPHATES BY MEANS OF INELASTIC SCATTERING OF COLD NEUTRONS*

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In hydrogen bond ferroelectrics of the KH_2PO_4 type, it is known that the hydrogen vibrations are fundamental to the ferroelectric phenomena as seen for example in the large shift in Curie point when hydrogen is replaced by deuterium.

Inelastic neutron scattering methods have proven successful in measuring the hydrogen vibrations in several hydrides.[1,2] Using this technique, measurements of the time-of-flight spectra of the scattered neutrons were made on three polycrystalline phosphate samples at room temperature and are shown in Fig. 1. In curve *A* for K_3PO_4, the high-energy peak can be attributed to PO_4 optical lattice vibrations. This is

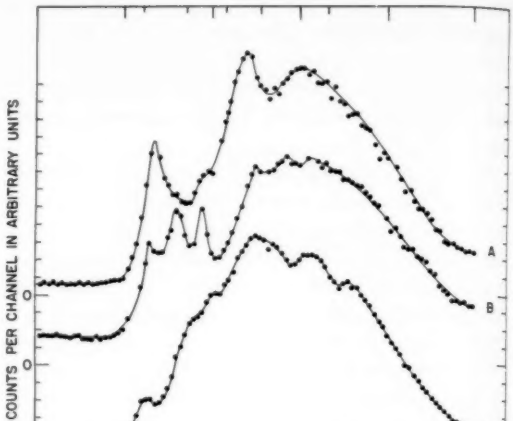

FIG. 1. Time-of-flight spectra of inelastically scattered cold neutrons (0.005 ev) from *A*, K_3PO_4; *B*, K_2HPO_4; *C*, KH_2PO_4.

confirmed by infrared measurements showing two reflection peaks at about 1100 cm^{-1} and 900 cm^{-1} for the PO_4 ion.[3] K_2HPO_4 spectra given in curve *B* show two added sharp peaks at 400 cm^{-1} and 550 cm^{-1}. These are certainly hydrogen vibration modes. From intensity consideration based on the population of vibration states at room temperature, one might argue that the 550-cm^{-1} peak represents the sum of the two transverse modes and that the 400-cm^{-1} peak is the longitudinal mode.

The ferroelectric KH_2PO_4 (paraelectric at room temperature) shows the high-energy PO_4 peak common to K_3PO_4 and K_2HPO_4. The hydrogen modes are merged into the low-energy region and there is an indication of peaking at 400 cm^{-1} which is also found in the K_2HPO_4. The KH_2PO_4 hydrogen vibration modes seem to be shifted to lower energies extending no higher than 550 cm^{-1}, and spread out over a larger energy region.

Our observation of low hydrogen vibration frequencies can be compared with infrared spectra measurements and values derived from neutron scattering density determinations. Many investigators have measured infrared spectra for deuterated and normal hydrogen bond ferroelectrics of the KH_2PO_4 type and looked for hydrogen vibration peaks. They have tried to identify them by examining shifts caused by the hydrogen-

deuterium mass difference. The wave numbers reported[4,5] for the light hydrogen are very high and range in value from 2800 cm^{-1} to 2400 cm^{-1}, in sharp contrast to our measurement. This energy range is not accessible to our energy gain experiment because these states will not be populated even near the dissociation temperature of the sample; however, the infrared peaks identified as hydrogen bond vibrations are assigned different frequency values by different investigators. They are always broad and very low in intensity. The complementary deuterium frequency is not always detected. The peaks do not sharpen in the ordered ferroelectric phase, in contrast to the other important peaks in the infrared spectrum. Therefore their assignments as hydrogen bond vibrations may be questionable. On the other hand, in KH_2PO_4 a sharp intense infrared reflectance peak at 540 cm^{-1} assigned to a combination mode with the lattice vibrations[3] lies at the high-energy end of the hydrogen vibration band observed in our measurement. This infrared peak becomes the most intense in the ordered ferroelectric phase, supporting our conclusion that it is a hydrogen vibration mode.

Bacon and Pease[6] observe in the paraelectric region an elongated shape for the scattering density of hydrogen, in the bond direction. They assume that it is a result of thermal anisotropy and thus obtain a zero-point energy of 0.08 ev and a low-level spacing of 0.16 ev. This corresponds to a wave number of 1250 cm^{-1}. We do not observe a peak at this energy, although it should be clearly measurable. The elongation therefore may result from a double potential well suggested as the other possibility in their paper.

Pirenne[7] in a theoretical treatment of KH_2PO_4 assumes a very low hydrogen vibration frequency of 148 cm^{-1} to account for the physical properties of the normal and deuterated compound. This energy region in our experiment has other modes of vibration that would make an assignment impossible.

Our observation of low hydrogen vibration frequencies will prove useful in understanding ferroelectric phenomena. The technique utilized in this work of studying fundamental lattice vibrations of correlated compounds can be employed to investigate even complex systems because of the direct interpretation that may be given to the data. It extends the range of observation of vibration modes much below that attainable by infrared technique.

We wish to thank W. Whittemore and General Atomic Division of General Dynamics for the use of their equipment in the preliminary stages of the work. We are indebted to D. J. Hughes and H. Palevsky for permission to run the samples at the Brookhaven Cold Neutron facility. Helpful discussions with B. C. Frazer, E. V. Sayre, G. H. Vineyard, and R. E. Weston are gratefully acknowledged.

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[1]Pelah, Eisenhauer, Hughes, and Palevsky, Phys. Rev. 108, 1091 (1957).

[2]Andresen, McReynolds, Nelkin, Rosenbluth, and Whittemore, Phys. Rev. 108, 1092 (1957).

[3]Murphy, Weiner, and Oberly, J. Chem. Phys. 22, 1322 (1954).

[4]Nakamoto, Margoshes, and Rundle, J. Am. Chem. Soc. 77, 6480 (1955).

[5]W. Kanzig, in Solid State Physics, edited by F. Seitz and D. Turnbull (Academic Press, New York, 1957), Vol. 4, p. 154.

[6]G. E. Bacon and R. S. Pease, Proc. Roy. Soc. (London) A230, 359 (1955).

[7]J. Pirene, Physica 15, 1019 (1949).

LOW-TEMPERATURE SPECIFIC HEATS OF BODY-CENTERED CUBIC Cr-Fe AND Cr-Mn ALLOYS*

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The electronic structure of the transition elements in the metallic state has not been fully understood as yet. It may be hoped that additional information with regard to the electronic specific heat coefficients and other physical properties of these elements, as well as of their alloys, will help in the interpretation of their structure.

On the basis of the band structure of the transition elements and of the usual interpretation of

the electronic specific heat coefficient, it should be possible to trace the shape of the partially filled d-band, i.e., the density of states as a function of the electron concentration, by determining experimentally the electronic specific heat coefficient γ for a series of solid solution alloys of a given crystal structure.

In the present work, low-temperature specific heats were determined for body-centered cubic binary solid solutions in the Cr-Fe and the Cr-Mn systems. Specific heats were measured in the temperature range of 1.6° to 4.2°K to an estimated accuracy of ±2% for most alloys. For a few alloys, where the specific heat depends very sensitively on the alloy composition, the accuracy may have been as low as ±5%. The alloy specimens were prepared by induction melting from electrolytically refined metals, and homogenized by annealing for at least three days at 1170°C in a purified 92% He + 8% H_2 gas mixture and quenched to room temperature.

The low-temperature specific heats were analyzed by means of the usual C_v/T vs T^2 extrapolation to $T=0$, to obtain the coefficient γ of the temperature-linear term. These γ values are plotted in Fig. 1 as a function of the electron concentration. The γ for Fe was measured and found to be in excellent agreement with the value reported by Keesom and Kurrelmeyer.[1] The γ for $Cr_{0.44}Fe_{0.56}$, as determined by Hoare and Matthews,[2] was plotted in Fig. 1, and it fits the curve for the Cr-Fe alloys very well. A large peak occurs at $Cr_{0.81}Fe_{0.19}$, with a γ value more than ten times that for Cr.[3] The fact that similarly high γ values were found also for specimens in the Cr-Mn alloy series at corresponding electron concentrations suggests that the observed effect may be related to characteristic values of the electron concentration. However, it is as yet uncertain whether or not the γ-peak should be interpreted simply in terms of density of states. The work on the low-temperature specific heat and on other physical properties of these alloys is continuing, and it should eventually make possible proper interpretation.

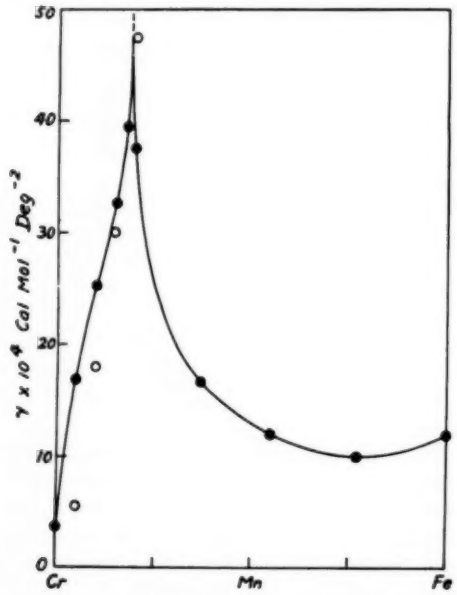

FIG. 1. Coefficient of term linear in temperature in the expression for the low-temperature specific heat, vs electron concentration. Cr-Fe alloys ●, Cr-Mn alloys ○.

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[1]W. H. Keesom and B. Kurrelmeyer, Physica 6, 663 (1939).

[2]F. E. Hoare and J. C. Matthews, Proc. Phys. Soc. (London) 71, 220 (1958).

[3]Friedberg, Estermann, and Goldman, Phys. Rev. 85, 375 (1952).

VALLEY-ORBIT SPLITTING OF ARSENIC DONOR GROUND STATE IN GERMANIUM

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A comparison of the density-of-states effective mass obtained from low-temperature Hall measurements on n-type germanium[1] with the effective-mass tensor yielded by cyclotron resonance experiments[2] suggests that, of the two states into which the donor ground state is split by the valley-orbit interaction,[3] the triplet is the lower one.[4] It has further been pointed out[5] that the differences in binding energies among the various donors is not inconsistent with such an assignment. By investigating the far-infrared excitation spectrum of bound electrons[6] in an elastically strained sample, we have now verified that for the arsenic donor the triplet is, in fact, the lowest lying state.

The splitting of the singlet, triplet, and a typical p-state under an applied strain can be calculated.[7] One assumes for this purpose that if the

wave function is expanded in conduction band Bloch waves only k-values very near the energy minima appear, so that the energy shift of each component is equal to the energy shift of the minimum near which it lies.

Since, for the p-state, the valley-orbit splitting is negligible, the quantization can be made simply by valleys. Such is not the case for the s-state, and as a result the frequencies of the s-p transitions are split. The intensities of the various components can be calculated from the selection rule that intervalley transitions by photon absorption are forbidden.

Figure 1 shows the splitting of any $(1s)$-(np) transition under a (111) uniaxial compression for the two cases in which the s-state is (a) a singlet or, (b) a triplet, taking into account the fact that the deformation potential Ξ_u is positive.[8] Relative intensities are indicated with each component. The undisplaced line in the triplet case arises from a nonzero electron population in substates other than the lowest one and is therefore temperature-sensitive. Except for this undisplaced line, the two patterns are mirror images of each other, with the intense line displaced toward high or low wavelength if the ground state is, respectively, the triplet or the singlet. The quantity Q, which governs the magnitude of the shifts, is given by

$$Q = \Xi_u t/9c_{44}, \tag{1}$$

where t is the applied stress and c_{44} the usual elastic constant.

We have studied the splitting of some $(1s)$-(np) absorption lines by such a uniaxial strain in a 1-mm thick sample containing about 3×10^{15} arsenic donors per cubic centimeter at a temperature of about 1.5°K. Figure 2 shows the observed absorption as a function of wavelength in the region of the $(1s)\rightarrow(2p, m=\pm1)$ transition, for the strained and unstrained case, after subtract-

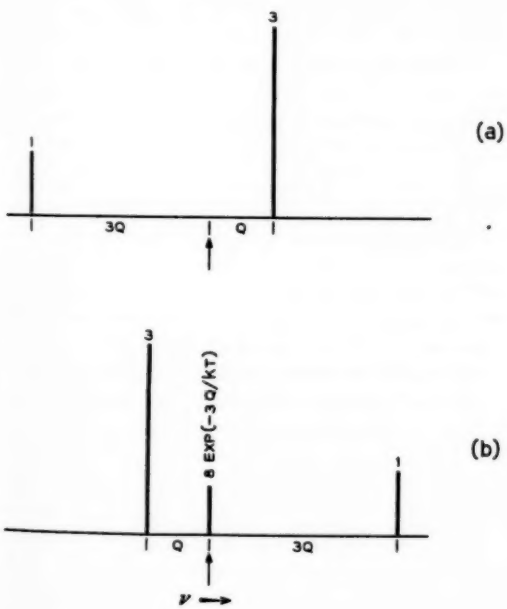

FIG. 1. Splitting of a $(1s\rightarrow np)$ line under a ⟨111⟩ compression. Case (a): singlet s-state. Case (b): triplet s-state.

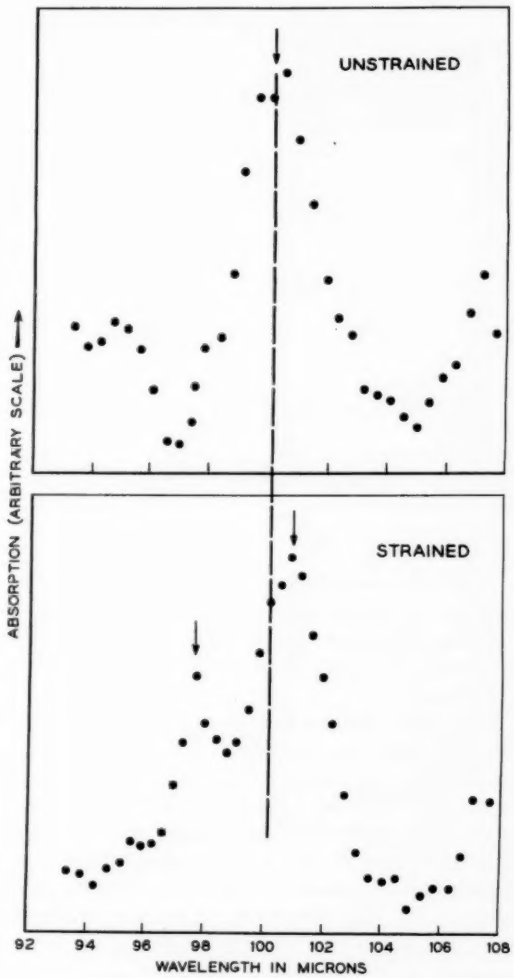

FIG. 2. Behavior of $(1s\rightarrow2p, \pm1)$ line under ⟨111⟩ compression; sloping background absorption subtracted out.

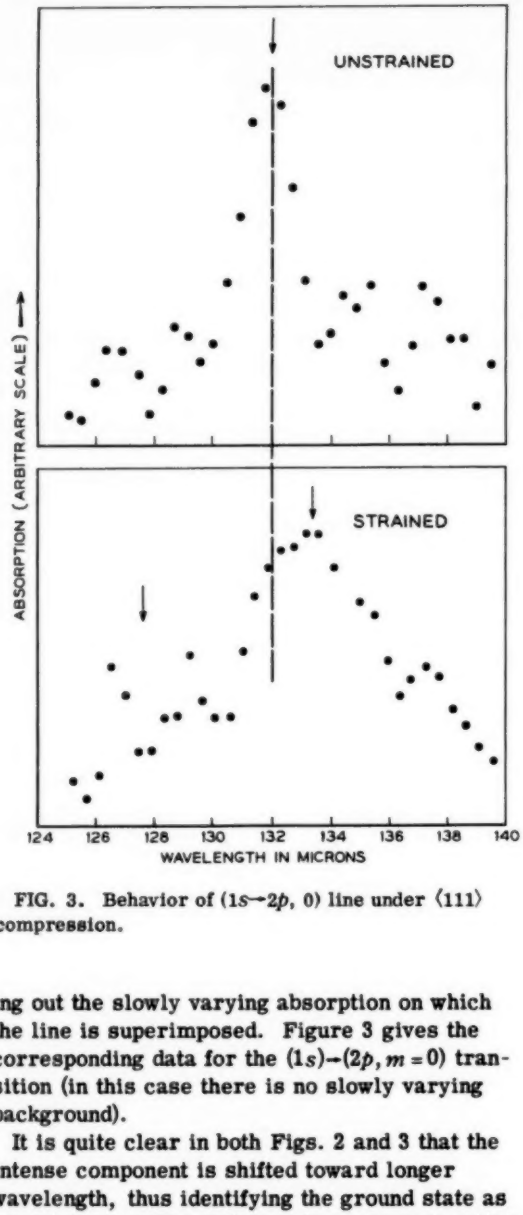

FIG. 3. Behavior of $(1s \rightarrow 2p, 0)$ line under $\langle 111 \rangle$ compression.

ing out the slowly varying absorption on which the line is superimposed. Figure 3 gives the corresponding data for the $(1s) \rightarrow (2p, m=0)$ transition (in this case there is no slowly varying background).

It is quite clear in both Figs. 2 and 3 that the intense component is shifted toward longer wavelength, thus identifying the ground state as the triplet. In Fig. 2 the weak component, which is shifted in the short-wavelength direction by triple the amount, is also apparent. This component is not, however, discernible in the 135μ region, Fig. 3, where it is masked by an oscillating background which we believe to be of instrumental origin. Since the $m=0$ transition has half the strength of the $m=\pm 1$ transition and is furthermore in a region where the power from the monochromator is smaller, it is not surprising that its weak component is more difficult to detect; we hope to improve this situation in future work. The undisplaced component, being quite weak at this temperature, cannot be definitely resolved in either set of data.

The arrows in Figs. 2 and 3 indicate a theoretical fit to the data. We treated the value of Q as an adjustable parameter and found it to be 0.85 cm^{-1}. (Of course, the same Q must apply to both the $m=0$ and $m=\pm 1$ transitions.) On the basis of Eq. (1), we would have predicted Q to be a little over 1 cm^{-1}. We do not at present consider this discrepancy significant, since our method of measuring the applied stress is quite crude.

We hope by extending the experiment to higher applied stress to make the strain-induced energy shift comparable to the triplet-singlet energy difference $\Delta\nu$; the results could then be interpreted to yield the value of $\Delta\nu$, whereas we have so far only determined its sign. We are also planning to repeat these measurements on donors other than arsenic.

[1]T. H. Geballe and F. J. Morin, Phys. Rev. 95, 1085 (1954).
[2]Lax, Zeiger, Dexter, and Rosenblum, Phys. Rev. 93, 1418 (1954).
[3]W. Kohn, in Solid State Physics, edited by F. Seitz and D. Turnbull (Academic Press, New York, 1957), Vol. 5.
[4]G. Weinreich, Bull. Am. Phys. Soc. Ser. II, 3, 102 (1958).
[5]G. Weinreich, Proceedings of the 1958 International Conference on Semiconductors (to be published).
[6]H. Y. Fan, Bull. Am. Phys. Soc. Ser. II, 3, 128 (1958).
[7]P. J. Price, Phys. Rev. 104, 1223 (1956).
[8]C. Herring and E. Vogt, Phys. Rev. 101, 944 (1956).

TIME-REVERSAL INVARIANCE AND PARITY CONSERVATION IN STRONG INTERACTIONS

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In view of the failure of parity conservation in weak interactions,[1] it is appropriate to investigate, under diverse situations, the degree to which strong interactions are invariant under the

operation of time reversal and space inversion.

One test of time-reversal invariance (TRI) can be derived from differential cross-section measurements of inverse reactions. A breakdown of the reciprocity property of the S matrix, which determines the transition probability per unit time, and hence of "detailed balance"[2] would indicate a lack of TRI. We have performed a "detailed balance" experiment with considerably higher accuracy than has heretofore been obtained[3] and under conditions which permit five competing open channels.

The experiment comprised a comparison of the differential cross sections for the reactions $p+t \rightleftarrows d+d$ at a center-of-mass energy of 6.3 Mev in the $p+t$ system. The $T(p,d)$ cross sections were measured directly,[4] whereas the $D(d,p)$ cross sections were synthesized from the energy dependence of the Legendre coefficients derived from previous measurements on the $D(d,p)$ reaction.[5]

Taking account of the consequences of the identity of the two emitted particles, in the forward reaction, on the available spin space and detection efficiency, a straightforward application of "detailed balance", predicts the ratio of emission cross sections to be $(\sigma_{p+t\rightarrow d+d})/(\sigma_{d+d\rightarrow t+p}) = 1.06$ in accord with the present experiment, the results of which are displayed in Fig. 1. According to recent calculations by Henley and Jacobsohn,[6] these results may be used to place an upper limit of ~2% on the fraction, F, of the matrix element which is odd with respect to TRI.

The second experiment to check time-reversal invariance, although independently conceived, is along lines recently proposed by Bell and Mandl[7] and Henley and Jacobsohn[6] and similar to one performed at much higher energy.[8]

The time-reversal experiment consists of two separate measurements. In one, a beam of protons, completely polarized in the direction of the z axis, and with velocity vector along the y axis, is scattered from a nucleus and the scattering asymmetry, A, measured in the x,y plane. (The experimental method and results for a number of elements have previously been described.[9]) A is defined as the quotient of the difference and sum of the intensities at points which are reflections by the y,z plane and is equivalent to the polarization, P, that would result from the scattering of an unpolarized beam of the same energy and at the same angle, provided the interaction is invariant under space and time inversion.[10] Nonequality of P and A, for nonzero spin targets, would signify the breakdown of time-reversal invariance if parity is conserved, whereas for spin-zero targets such nonequality would require parity nonconservation. The second measurement therefore entails a direct determination of P, He^4 being used as the analyzer.

A and P have been compared for Co, Fe, and Ni at an energy of 7.1 Mev and at an angle of 96°. Table I presents the results, which are seen to be consistent with TRI and parity conservation. Ni and Fe are sensitive to parity plus time reversal while Co is sensitive to time reversal irrespective of whether parity is conserved.

The second experiment which we performed to

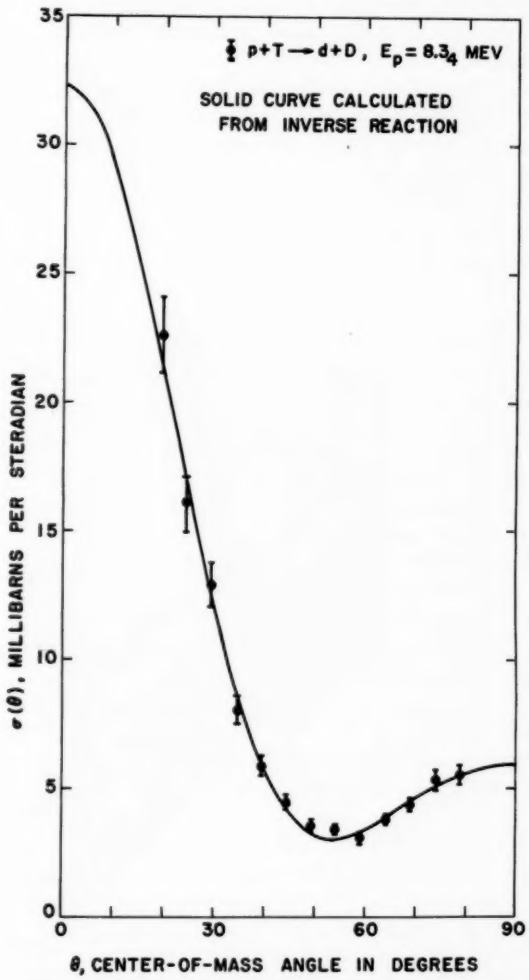

FIG. 1. Comparison of the experimentally observed values of the differential cross section for the $T(p,d)D$ reaction with those calculated from the measurements on the inverse reaction.

Table I. Comparison of asymmetry and polarization for Ni, Fe, and Co.

Element	P	A
Ni	0.13±0.02	0.14±0.03
Fe	0.11±0.03	0.16±0.04
Co	0.19±0.03	0.16±0.03

test parity conservation is much more sensitive and parallels similar experiments at high energies.[11] When 10-Mev polarized protons are incident on complex nuclei, large azimuthal asymmetries are observed in the x,y plane, as shown in Fig. 2. (Positive A implies a preponderance of protons twice scattered to the left, once by the polarizer and once by the target, over those scattered once to the left and once to the right.) Following the suggestion of Lee and Yang,[12] we have compared scattered intensities at points in the y,z plane which are reflections by the x,y plane. If F denotes the ratio of the fractional strength of the interactions which mix parities, the ratio of intensities in the direction along and opposite to the polarization vectors are related by $R \equiv [(1+F)/(1-F)]^2$. We measured R for a number of elements at various angles including those corresponding to the absolute maxima and minima in Fig. 2. The data are presented in Table II. Our results place an upper limit of 10^{-3} on F^2. This may be taken as confirmation, from still another direction, of results based on parity-forbidden reactions and nonappearance of circular polarization in nuclear gamma-ray transitions.[13]

The authors are indebted to the cyclotron group for their best efforts in making available intense beams of protons and alpha particles of the appropriate energies.

It is a pleasure to acknowledge the theoretical guidance of Dr. R. M. Thaler in all phases of this work. We wish to thank Dr. Leona Marshall for enlightening discussions on time reversal invariance.

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[1]Wu, Ambler, Hayward, Hoppes, and Hudson, Phys. Rev. 105, 1413 (1957); Garwin, Lederman, and Weinrich, Phys. Rev. 105, 1415 (1957).

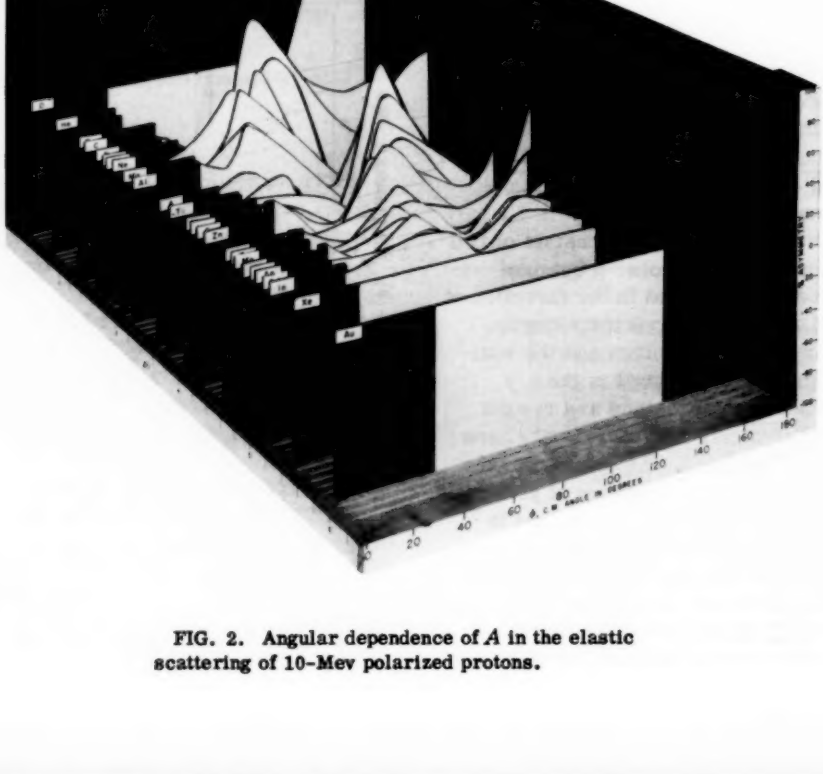

FIG. 2. Angular dependence of A in the elastic scattering of 10-Mev polarized protons.

Table II. Ratio of intensities parallel and antiparallel to the polarization vector. The standard error on each point is ± 4%.

Element / Angle	He^4	C	N	O	Ne	Al	Fe	Cu	Ni
50°			1.03	1.05	0.98				
60°						0.97			
70°		1.06					1.05	0.96	0.97
80°			0.93	1.00	0.97				
90°	1.04	0.98							
100°						1.04	0.97	1.03	0.98
110°		0.96	0.94	1.04	1.03				
130°	0.95								

[2]J. M. Blatt and V. F. Weisskopf, Theoretical Nuclear Physics (John Wiley and Sons, New York, 1952), p. 530; E. M. Henley and B. A. Jacobsohn, Phys. Rev. 108, 502 (1957); E. Fermi, Nuclear Physics (The University of Chicago Press, Chicago, 1950).

[3]J. H. Coon, Phys. Rev. 80, 488 (1950); Willard, Bair, and Kingston, Phys. Rev. 90, 865 (1953); W. Selove, Phys. Rev. 103, 136 (1956); J. E. Perry, Jr., and S. J. Bame, Jr., Phys. Rev. 99, 1368 (1955); E. G. Fuller, Phys. Rev. 96, 1306 (1954); J. B. Reynolds and K. G. Standing, Phys. Rev. 101, 158 (1956); Levine, Bender, and McGruer, Phys. Rev. 97, 1249 (1955); Bodansky, Eccles, Farwell, Rickey, and Robison, Bull. Am. Phys. Soc. Ser. II, 3, 327 (1958); Burcham, McCauley, Bredin, Gibson, Rotblat, and Prowse, Nuclear Phys. 5, 141 (1958).

[4]J. N. Bradbury and L. Stewart, Bull. Am. Phys. Soc. Ser. II, 3, 417 (1958).

[5]Brolley, Putnam, and Rosen, Phys. Rev. 107, 820 (1957).

[6]E. M. Henley and B. A. Jacobsohn, Phys. Rev. (to be published).

[7]J. S. Bell and F. Mandl, Proc. Phys. Soc. (London) 71, 272 (1958).

[8]Heer, Roberts, and Tinlot, Phys. Rev. 111, 645 (1958).

[9]L. Rosen and J. E. Brolley, Jr., Phys. Rev. 107, 1454 (1957); L. Rosen and J. E. Brolley, Jr., Proceedings of the Second United Nations International Conference on the Peaceful Uses of Atomic Energy, Geneva, 1958 (United Nations, Geneva, 1958), Paper UN-668.

[10]L. Wolfenstein, Annual Review of Nuclear Science, (Annual Reviews, Inc., Palo Alto, 1956), Vol. 6, p. 43.

[11]Oxley, Cartwright, and Rouvina, Phys. Rev. 93, 806 (1954); Chamberlain, Segrè, Tripp, Wiegand, and Ypsilantis, Phys. Rev. 93, 1430 (1954).

[12]T. D. Lee and C. N. Yang, Phys. Rev. 104, 254 (1956).

[13]N. Tanner, Phys. Rev. 107, 1203 (1957); D. H. Wilkinson, Phys. Rev. 109, 1603 (1958).

TIME-REVERSAL INVARIANCE AND THE INVERSE REACTIONS $C^{12}+\alpha \rightleftharpoons N^{14}+d$*

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It was pointed out by Henley and Jacobsohn[1] that although time-reversal invariance in strong interactions can be investigated by examining detailed balance in an appropriate nuclear reaction and its inverse, no very precise comparisons of this nature were then available.[2] Subsequently, relatively precise polarization and asymmetry measurements have indicated that for p-p scattering at energies near 200 Mev the magnitude of the time-reversal noninvariant term of the scattering matrix is not more than a few percent of the average magnitude of invariant terms.[3] The experiment described here approaches the problem through the study of detailed balance in nuclear reactions. It further differs from the p-p experiments in that it also involves n-p forces and is in a quite different energy region.

Following the suggestion of Henley and Jacobsohn, we have measured the angular distributions for the inverse (ground state) reactions $C^{12}(\alpha,d)N^{14}$ and $N^{14}(d,\alpha)C^{12}$. The measurements were carried out using 41.7-Mev alpha particles and 20.0-Mev deuterons from the external beam of the University of Washington cyclotron. (It was necessary to degrade the normal 21-Mev deuteron beam in a thin absorber in order to

match correctly the natural 41.7-Mev alpha energy.) The energies of the beams were determined from measurements of the range of particles in aluminum, using recent range-energy data.[4] Carbon and nitrogen targets were obtained by using thin foils of polystyrene and Cymel,[5] respectively.

The deuterons in the (α, d) measurement were detected in two NaI(Tl) or CsI(Tl) scintillation counters (used in a "$dE/dx-E$"-type arrangement), while the alpha particles in the (d, α) measurement were detected in a counter combination consisting of a xenon ionization chamber and a scintillation counter. The essentially monoenergetic particle group of interest was identified by demanding a coincidence between pulses of proper magnitudes from the two counters of a detector. This was accomplished by displaying pulses from both counters on twenty-channel analyzers and establishing coincidences between events in selected channels of the two analyzers.

The observed angular distributions, converted to the center-of-mass system, are shown in Fig. 1. As no attempt was made to establish absolute cross sections, the distributions are arbitrarily normalized to give the closest over-all agreement. The two angular distributions show striking qualitative agreement, and more than half the points for each reaction lie within one probable error (typically about 4%) of the common smooth curve.

However, as can be seen in Fig. 1 and in the analysis by angular intervals given in Table I, consistent differences between the two distributions appear in certain regions, notably near 90° and near 130°. A re-examination of experimental difficulties suggests that the differences near 90° may be due to an underestimate of background contributions to the (α, d) counting rates and that the difference near 130° may be due to mechanical errors in determining the (d, α) angles. (Contributions to the errors from other sources, such as incorrect beam energies and multiple scattering losses, are believed to be comparatively small.) These differences are not outside the estimated maximum possible error, and we conclude that they do not represent significant evidence for a departure from detailed balance.

In an attempt to obtain an over-all characterization of the degree of agreement between the

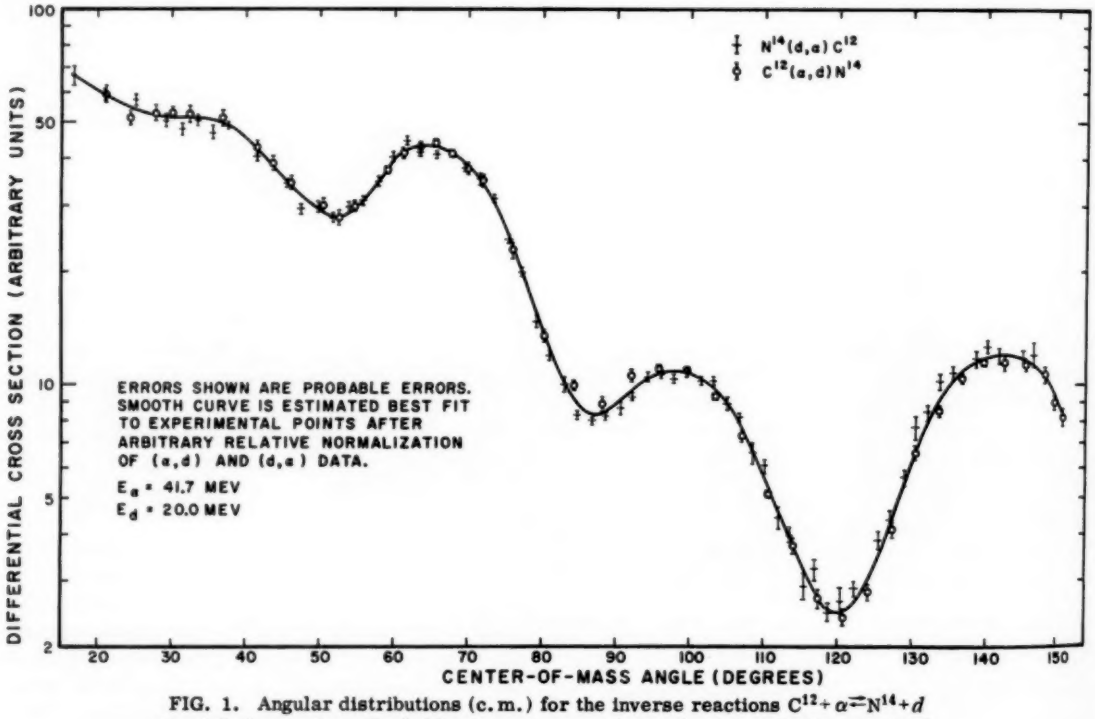

FIG. 1. Angular distributions (c.m.) for the inverse reactions $C^{12}+\alpha \rightleftharpoons N^{14}+d$ at matched energies. Probable errors on experimental points include statistical errors and estimated uncertainties in the subtraction of small background contributions to the measured counting rates.

Table I. Analysis of differences between distributions. For each 10° interval a deviation, d, from the smooth curve of Fig. 1 is found by (a) determining the fractional discrepancy between each experimental point and the smooth curve, (b) taking the algebraic sum of these discrepancies, and (c) dividing by the number of points in the interval. The difference, D, between the two angular distributions in any 10° interval is the algebraic difference between the deviations, d. In computing the average of the magnitude of D, noted at the bottom of the table, the smallest difference is omitted to compensate for the arbitrary relative normalization of the two distributions. For qualitative orientation, a typical probable error, ϵ, for individual points within each interval is also tabulated.

Angular interval (c.m. degrees)	$C^{12}(\alpha,d)N^{14}$ d (percent)	ϵ (percent)	$N^{14}(d,\alpha)C^{12}$ d (percent)	ϵ (percent)	Difference D (percent)
20-30	-1.2	5.0	+1.4	4.3	-2.6
30-40	+2.7	4.4	-3.6	3.4	+6.3
40-50	+1.6	4.0	-2.7	2.8	+4.3
50-60	+1.0	3.5	+1.0	2.7	0.0
60-70	-0.6	2.7	-1.8	2.2	+1.2
70-80	+1.1	4.3	+0.1	2.9	+1.0
80-90	+4.5	3.8	-2.9	2.8	+7.4
90-100	+4.7	3.1	-2.1	2.6	+6.8
100-110	-4.4	3.2	+1.3	3.6	-5.7
110-120	-1.4	4.2	+2.2	6.4	-3.6
120-130	-6.3	5.0	+5.5	6.6	-11.8
130-140	-4.3	2.9	+7.3	5.0	-11.6
140-150	-1.7	4.3	+3.6	4.6	-5.3
				Average difference: $\langle \lvert D \rvert \rangle_{Av}$	=5.6%

two distributions, we have calculated an average difference in the manner described in Table I. As seen from Table I, the calculated average difference between the measured angular distributions is less than 6% and would be smaller if the somewhat less accurate data beyond 110° were given reduced weight. Differences of this order can be understood in terms of experimental errors, and we believe that the present data, considered as a whole, are not inconsistent with the assumption of exact detailed balance. While experimental errors could tend to cancel an actual difference between the distributions over a limited angular region, such accidental reduction is unlikely over a wide region. Therefore (aside from a single unmeasured normalization factor) it is concluded that the differential cross sections for the reactions $N^{14}(d,\alpha)C^{12}$ and $C^{12}(\alpha, d)N^{14}$ from 20° to 150° (c.m.) probably differ on the average by less than 6%. This agreement is interesting in view of the wide, and somewhat irregular, fluctuations of the cross sections with angle.

A precise assessment of the upper limit placed by these results upon the fraction of the Hamiltonian which is noninvariant with respect to a time inversion must await analysis in terms of a specific model which can account satisfactorily for the observed angular distributions and which in addition can predict their sensitivity to possible violations of time reversal invariance. In the absence of such a theory, we apply the approximate criterion of Henley and Jacobsohn,[1] and conclude that the present results imply that this noninvariant fraction is probably less than 3%.

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[1]E. M. Henley and B. A. Jacobsohn, Phys. Rev. 108, 502 (1957); Phys. Rev. (to be published). We are indebted to Dr. Henley and Dr. Jacobsohn for many helpful discussions.

[2]Work reported recently includes Burcham, McCauley, Bredin, Gibson, Prowse and Rotblat, Nuclear Phys. 5, 141 (1958); Bodansky, Eccles, Farwell, Rickey, and Robison, Bull. Am. Phys. Soc. Ser. II, 3, 327 (1958); J. N. Bradbury and L. Stewart, Bull Am. Phys. Soc. Ser. II, 3, 417 (1958); L. Rosen (private communication to B. A. Jacobsohn). See reference 1 for

bibliography of earlier work.

[3]Hillman, Johansson, and Tibell, Phys. Rev. 110, 1218 (1958); and A. Abashian and E. M. Hafner, Phys. Rev. Lett. 1, 255 (1958).

[4]H. Bichsel (private communication).

[5]We are grateful to Dr. J. K. Magrane and Mr. H. R. Russell of the American Cyanamid Company for their assistance in providing us with foils of this nitrogen compound.

NEUTRON SPECTRA FROM $Be^9(He^3,n)C^{11}$ AND $Be^9(He^4,n)C^{12}$ USING PULSE SHAPE DISCRIMINATION IN AN ORGANIC SCINTILLATOR

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It has been shown recently[1] that it is experimentally possible to use pulse shape discrimination to differentiate between neutron and gamma-ray induced pulses from some organic scintillators. This provides a powerful method for the study of fast-neutron spectra from nuclear reactions and also for the study of gamma-ray spectra, from a NaI(Tl) crystal, in coincidence with fast neutrons.

The purpose of this Letter is to report some measurements on the suitability for pulse shape discrimination of some commercially available organic scintillators and to report some preliminary observations of the high-energy neutron groups from the $Be^9(He^4,n)C^{12}$ and the $Be^9(He^3,n)C^{11}$ reactions.

The pulse shape discrimination method is based upon the fact that the ratio of the intensity of the long-period decay component (~200 mμsec) to that of the short-period component (~4mμsec) is different for recoil protons and electrons.[1] This difference was assessed for a number of scintillators by taking two pulses from the photomultiplier, a current pulse from the anode, and an integrated pulse from a lower dynode. After suitable amplification and stretching, a narrow band of current pulses was selected by a single-channel analyzer, the output of which opened a linear gate to allow the corresponding integrated pulses to be analyzed by a 100-channel pulse amplitude analyzer. If the pulses from the organic scintillator are caused by the neutrons and gamma-rays from a PuBe source two peaks will usually appear on the kicksorter, the higher being composed entirely of neutron pulses and the lower of gamma-ray pulses. A gamma-ray source produces a single peak. Table I lists the ratio, R, of the channel number of the neutron peak to the channel number of the gamma-ray peak (the corresponding electron energy is approximately 1 Mev), for a number of commercial organic scintillators.[2,3]

As part of a program to study neutron spectra and (n,γ) coincidences using a neutron detector which exhibits neutron gamma-ray discrimination, we have started a series of measurements with the liquid scintillator *NE*212.[2] The reactions $Be^9(He^3,n)C^{11}$ and $Be^9(He^4,n)C^{12}$ were chosen for study. The former reaction, which has not been studied previously, yields a number of neutron groups together with the associated gamma-rays, and gamma-rays from the reaction $Be^9(He^3,p\gamma)B^{11}$ are also emitted.[4] By using a mixture of He^3 and He^4 in the ion source of the Chalk River electrostatic accelerator the reactions could be studied by simply changing the deflecting-magnet field. The scintillator was placed at 0° and at a distance of 3 in. from a thin Be^9 target which was then bombarded by either 1.9-Mev He^3 or He^4 ions. The neutron output was monitored with a ZnS(Ag)-loaded plastic scintillator, *NE*404,[2] placed at 90° to the beam direction.

The separation of the neutron pulses and the gamma-ray pulses, from the *NE*212 scintillator, over a broad range of neutron energies, was effected by a circuit similar to that devised by Owen.[1] In this circuit the voltage between the last dynode and anode of a photomultiplier (Dumont 6363) was reduced to about 4 volts so that the current, leaving the last dynode during the "fast" part of the current pulse, was space-charge limited. The signal from the last dynode was consequently negative during the "fast" part

Table I. Neutron–gamma-ray discrimination properties of some organic scintillators. R is the ratio of the channel number of the neutron peak to that of the gamma-ray peak.

Scintillator	R
*NE*202	1.17
*NE*210	1.11
*NE*211	1.11
*NE*212	1.30
Stilbene	1.33

of the current pulse but became positive during the "slow" part to an extent depending upon the relative intensity of the slow component to the fast component. The output of a pulse-height discriminator, set on the amplified positive pulse from the last dynode, was then used to open a gate for a linear pulse obtained from an earlier dynode. The output of the linear gate was displayed on a 100-channel pulse amplitude analyzer. Owen has used this circuit to display the neutron and gamma-ray pulse spectrum from a PuBe source.[1]

Figure 1 shows the pulse spectra obtained from the $Be^9 + He^3$ reaction for three values of the discriminator bias on the pulses from the last dynode. As expected, the effect of increasing the bias was first to eliminate the gamma-ray pulses at approximately bias 700 and then to cut off greater portions of the low-energy proton recoil spectrum. The bias 800 curve represents only the proton recoil distribution above 3 Mev proton energy. Experiments with PuBe neutrons and the 2.61-Mev gamma-rays from a radiothorium source (the latter give pulses corresponding to neutron energies of up to approximately 5 Mev) showed that >99% of the gamma-rays were eliminated from the PuBe neutron spectrum. In Fig. 2 the spectra are plotted on a logarithmic scale to show the effect of discrimination more clearly. Also shown is a calibration recoil proton pulse spectrum from the $Be^9(He^4,n)C^{12}$ reaction taken under the same conditions (bias 800) as the spectrum from $Be^9(He^3,n)C^{11}$. The steps in the $Be^9(He^3,n)C^{11}$ spectrum at approximately 9.5, 7.6, and 5.0 Mev correspond to neutron groups feeding the ground, first excited, and unresolved second and third excited states, respectively.[5] The step at 5.0 Mev is broader than expected from a single neutron group.

It is intended to pursue these and related measurements further, using larger tanks of liquid scintillator[6] and other methods of separating the neutron and gamma-ray pulses.

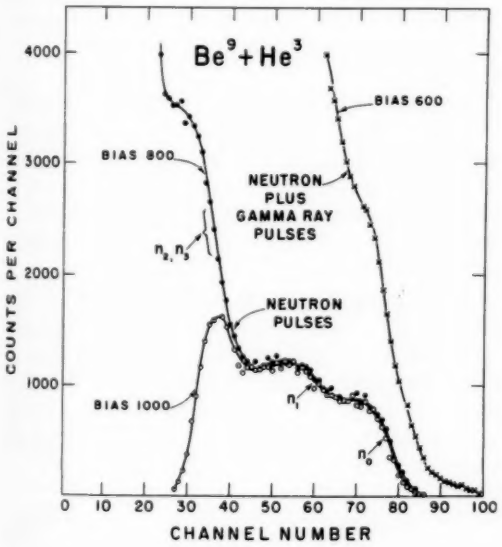

FIG. 1. Pulse-height spectra, obtained with a liquid scintillator, from the $Be^9 + He^3$ reaction, showing the effect of pulse shape discrimination. The steps n_0, n_1, n_2, and n_3 correspond to neutron groups feeding the ground, first excited, and unresolved second and third excited states respectively in C^{11}.

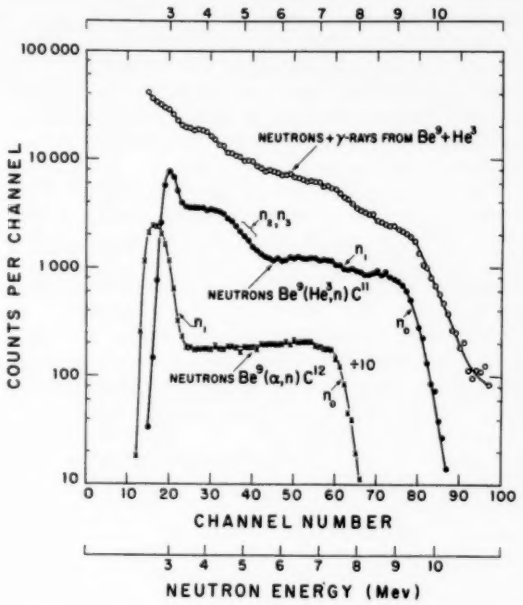

FIG. 2. Pulse-height spectra from the $Be^9(\alpha,n)C^{12}$ and $Be^9(He^3,n)C^{11}$ reactions.

[*]Seconded from the United Kingdom Atomic Energy Authority.

[1]R. B. Owen, Nucleonics 16, No. 6, 54 (1958); IRE Trans. on Nuclear Sci. NS-5, 198 (1958); and a paper to be published in Onde Electrique. The authors are indebted to R. B. Owen for the communication of his results prior to publication. It appears that the pulse-shape discrimination method was first suggested by F. D. Brooks, Progress in Nuclear Physics (Butterworths-Springer, London, 1956), Vol. 5, p. 284. An

experimental study of the effect together with a theoretical discussion has been published by G. P. Wright, Proc. Phys. Soc. (London) B69, 358 (1956).

[2]The encapsulated liquid scintillators, 1½-in. diam by 1½ in. long and the fast-neutron detector *NE*404, were obtained from Nuclear Enterprises, Winnipeg, Manitoba, Canada.

[3]The authors wish to acknowledge the cooperation of Professor L. Katz during early work on some of these scintillators and during unpublished work on a variety of inorganic scintillators.

[4]Ferguson, Gove, Kuehner, Litherland, Almqvist, and Bromley, Phys. Rev. Lett. 1, 414 (1958).

[5]V. Johnson, Phys. Rev. 86, 302 (1952).

[6]Preliminary observations, of pulse-height spectra from a 3-in. diam by 3-in. long *NE*212 liquid scintillator, indicate that the shape of the high-energy portion of the spectrum of pulses from monoenergetic neutrons (7-10 Mev) is less favorable for the analysis of complicated neutron spectrum than the spectrum from a smaller scintillator. See J. E. Hardy, Rev. Sci. Instr. 29, 705 (1958).

p-p S-WAVE PHASE SHIFTS FROM A BOUNDARY CONDITION POTENTIAL*

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Nucleon-nucleon scattering experiments have led to a potential model for the two-nucleon interaction that conventionally features a very hard repulsive core followed by some type of meson-theoretic attractive well.[1,2] While the hard repulsive core is needed to explain very high-energy nucleon-nucleon scattering, Moszkowski and Scott have pointed out recently[3] that the possibility exists of explaining lower energy (say below 200 Mev) nucleon-nucleon scattering by means of a potential featuring a somewhat softer core. The motivation for their remark is that an infinite repulsive core leads to difficulties when the many-body problem is expressed in terms of two-nucleon scattering amplitudes. If a modified potential exists that can give agreement with nucleon-nucleon scattering over the energy range of importance—with the core region chosen so as to eliminate the divergences that result from conventional perturbation expansions—then this model would be a convenient starting point for many-body calculations.[4]

(In the discussion of the present Letter, we restrict ourselves to the singlet-*S* nucleon-nucleon scattering state. *S*-state scattering is more sensitive to details of the core region than are scatterings from the higher angular momentum states. Hence if satisfactory results are obtained for the *S*-state, then it is reasonable to expect that a modified core can also be made to give agreement for *P* and higher waves. Also, by treating only the *S*-state, we side-step the complexities associated with tensor and/or spin-orbit terms that are necessary for a complete nuclear potential model, but which do not essentially relate to a modification of the central core potential.)

Moszkowski and Scott[3] chose a core potential that can be represented by assigning to the wave function at the core radius c the boundary condition

$$\left(\frac{1}{u}\frac{du}{dr}\right)_{r=c} = \frac{1}{c}. \tag{1}$$

This boundary condition has the property that at zero energy the two-nucleon wave function goes over into the free-particle wave function for separations $r>c$. Mathematically, the core consists of an infinitely-repulsive region followed by a region that is infinitely attractive and of infinitesimal width, chosen so that the effects of the attractive and repulsive parts cancel out at zero energy. This is a convenient potential to use from the standpoint of Brueckner theory applications.[4]

The question remains as to whether a two-nucleon potential featuring the boundary condition given by Eq. (1) can yield the same results as a hard core potential.[1,2] Moszkowski and Scott[3] showed that for the simple case of a core followed by a square well, the two models could be made to agree to within 1/2 degree for the singlet-*S* phase shift over the laboratory energy range 0-250 Mev. This was an analytic result. Calculations for the more realistic case of a core followed by a Yukawa well require a computing machine. Such calculations were carried out on the Livermore UNIVAC, using conventional iteration techniques. If we define the potential as

$$V = -V_0\, e^{-\mu r}/\mu r, \quad r>c \tag{2}$$

and use the boundary condition from Eq. (1), then we obtain the results shown in Table I, which are there compared with the model of Gammel and Thaler[1] and with the model of Signell and Marshak.[2] Four sets of parameters for the Moszkowski-Scott potential are shown. The first two

Table I. Singlet-S phase shifts from a Yukawa model with a core.

Model	Core	Core radius (in 10^{-13} cm)	μ[a] (in 10^{13} cm^{-1})	V_0[a] (Mev)	Singlet-S phase shift (deg) 10 Mev	40 Mev	150 Mev
Gammel-Thaler	∞	0.4	1.45	425.5	54.3	39.8	11.7
Signell-Marshak[b]					(50)	42.9	19.2
Moszkowski-Scott	Eq. (1)	0.765	1.55	344	53.3	41.1	11.7
Moszkowski-Scott	Eq. (1)	0.82	1.55	395	55.9	41.3	10.5
Moszkowski-Scott	Eq. (1)	0.74	1.50	320	52.5	45.8	20.5
Moszkowski-Scott	Eq. (1)	0.74	1.55	350	51.2	42.5	14.4

[a]See Eq. (2). [b]Gartenhaus potential.

sets are in good agreement with the Gammel-Thaler phase shifts. The last two sets show that a fit to the Signell-Marshak phase shifts is also possible, although no attempt was made to find a set of parameters giving precise agreement with these. (The Signell-Marshak phase shifts are actually n-p phase shifts and should not be directly compared with the other calculations, which are p-p phase shifts; however, the comparison is still of some interest.)

It can be seen from Table I that when the boundary condition from Eq. (1) is used, the core radius must be about double the value used by Gammel and Thaler. This was also true for the square well calculation of Moszkowski and Scott.[3] Representing the core by a boundary condition such as Eq. (1) means that we are only interested in the asymptotic form of the wave function as it emerges from the core, and the radius at which we define the logarithmic derivative to be Eq. (1) is in some sense a correlation radius.

The conclusion we draw from the present work and from the paper by Moszkowski and Scott[3] is that the boundary condition model specified by Eq. (1) and a standard external potential is as adequate as the hard-core models in serving as a representation of low-energy (say 0 to 150 Mev) nucleon-nucleon scattering data.

The present calculations were carried out at the suggestion of Professor Moszkowski.

*This work was performed under the auspices of the U. S. Atomic Energy Commission.

[1]J. L. Gammel and R. M. Thaler, Phys. Rev. 107, 291 (1957).

[2]P. S. Signell and R. E. Marshak, Phys. Rev. 109, 1229 (1958).

[3]S. A. Moszkowski and B. L. Scott, Phys. Rev. Lett. 1, 298 (1958).

[4]S. A. Moszkowski and B. L. Scott (to be published).

HELICITY OF NEGATIVE MUONS FROM PION DECAY*

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(Received December 30, 1958)

A method for the direct measurement of the helicity of μ^- mesons from pion decay has been suggested by Jackson, Treiman, and Wyld.[1] We now report the preliminary results of such an experiment, which involves the detection of the asymmetry of β-decay of B^{12} produced by polarized muons via the reaction $\mu^- + C^{12} \rightarrow B^{12} + \nu$.

The experimental arrangement is shown in Fig. 1. The muon beam was stopped in C_5H_{12}

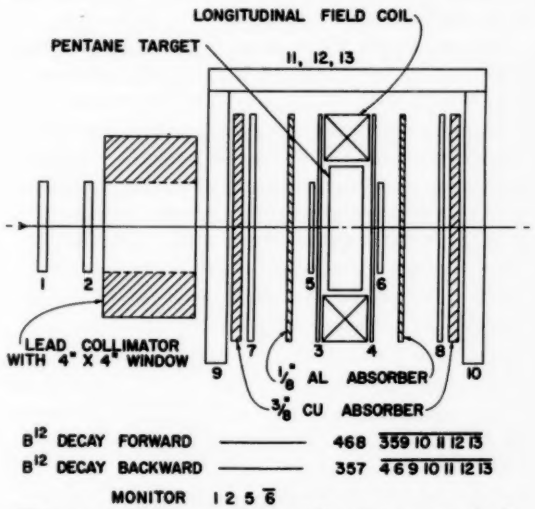

FIG. 1. Counter arrangement for measurement of B^{12} asymmetry. Beam is incident from the left. Not shown are the Al beam absorber and the vertical field coil.

contained in a Mylar-walled $5\frac{1}{2}$ in. $\times 5\frac{1}{2}$ in. $\times 1\frac{1}{2}$ in. target. All muons stopped in the pentane were captured in carbon atoms, where $1\frac{1}{2}$ to 2% underwent nuclear capture to bound states of B^{12}.[2,3] The β-rays resulting from the decay of B^{12} ($T_{1/2} = 20.34 \pm 0.5$ msec)[4] were detected by a pair of three-counter telescopes which we denote F and B for forward and backward, respectively. Various anticoincidence counters shielded against background, which was largely cosmic radiation. The following timing sequence was used (see Fig. 2). The cyclotron was turned on for eight

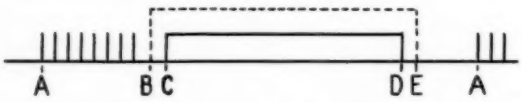

FIG. 2. Timing sequence. A is the first of eight cyclotron FM pulses, B the start of the longitudinal field pulse (condition U), C the start of the 90-msec counting period, D the close of the 90-msec counting period, and E the end of the longitudinal field pulse.

FM pulses, which occupied 40 msec, and was then turned off for 130 msec. Delayed coincidences occurring in each telescope during a 90-msec gate which started 13.6 msec after the last preceding FM pulse were scaled and also timed relative to that pulse with a time interval meter. The 13.6-msec delay was chosen to reduce the contribution to the counting rate of a cyclotron-produced activity with a half-life of approximately 4 msec.

Measurements were made under two conditions: condition P, in which the target was subjected to a dc field of 100 gauss parallel to the beam direction and condition U, in which this field was applied only during 102-msec intervals encompassing each 90-msec counting period. At all times a transverse dc field of 9 gauss was applied by a rectangular coil wound on the target. Thus, counts were registered only when both fields were present in order to avoid systematic errors. In addition, each scintillation detector was magnetically shielded, and found to be insensitive to the applied fields. In condition P the resultant magnetic field was such as to preserve the longitudinal polarization which the B^{12} received from the muons. In condition U the transverse field served to depolarize the B^{12} before the longitudinal field was turned on and the counting period began. P and U counts corresponding to a fixed monitor count were taken in an ordered sequence so as to minimize the effect of electronic drifts. The asymmetry, defined as $\epsilon + 2[(P-U)/(P+U)]$ was computed for each telescope from the scaler counts.

The negative muons were produced in the Carnegie Tech cyclotron and had an initial energy of 115 Mev as measured by range in Al. We have made measurements with a stopped μ^+ beam of similar energy, and found positrons emitted preferentially backwards. From the results of asymmetry measurements on the π-μ-e chain from pions decaying at rest,[5] and also from consideration of the available proton energy in the cyclotron, we conclude that the muons in the μ^+ beam, and also the μ^- beam, are moving forward relative to their initial directions in the pion system. In order to determine that the stopped muons were polarized, a counter telescope in the forward direction was used to detect the decay electrons from muons stopping in a cylindrical pentane target. The time of emission of each electron relative to the stopping muon was measured with a time interval meter to an accuracy of 10^{-7} sec. A 40-gauss transverse field along the target axis caused precession of the muons with a 1.75-μsec period, which was readily observed as a sinusoidal modulation of the decay curve. An analysis of 1.5×10^4 counts indicates an initial backward excess of electrons and about 15% muon polarization.

The decay of the B^{12} formed when muons were stopped in the rectangular pentane target is shown in Fig. 3. An analysis of 2.7×10^4 counts with the aid of an IBM 704 computer gives a half-life for the exponential component of 20.4 ± 1.1 msec, in good agreement with recent measurements on B^{12}.

Pions were stopped in the target, and a time analysis of the activities produced was obtained in the same manner as for muons with the length of the gate increased to 250 msec and the time between cyclotron pulses increased accordingly. Results are shown in Fig. 4 and indicate a two-component decay, one of about 20 msec half-life which we attribute to B^{12} and the other of much longer lifetime. A further effort to identify the long component was made by scanning over a 1.04-sec period. A half-life of 0.8 sec was obtained which may be Li^8 or B^8. A more precise determination of the nature of this long-lived activity awaits measurements with an extended counting gate. However, we find no evidence for the 80-msec activity reported by the CERN group.[3]

With muons stopping in the target we obtained a

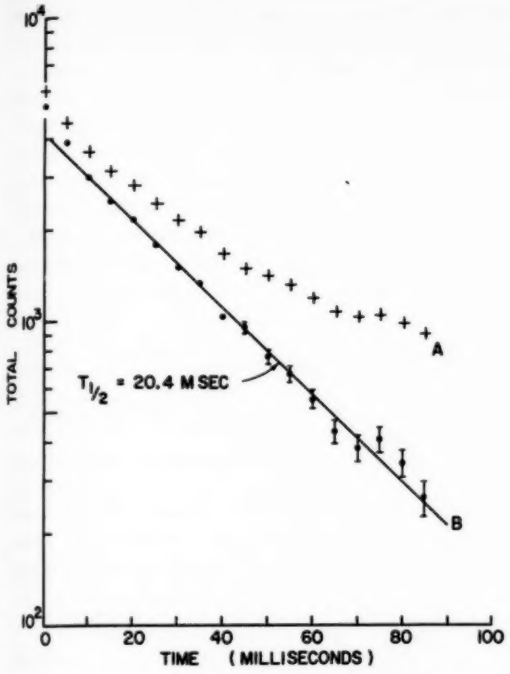

FIG. 3. Curve A is the decay curve observed with muons stopping in the target. Curve B is the result of subtracting the computed background. Each point represents the total number of counts obtained in a 5-msec interval.

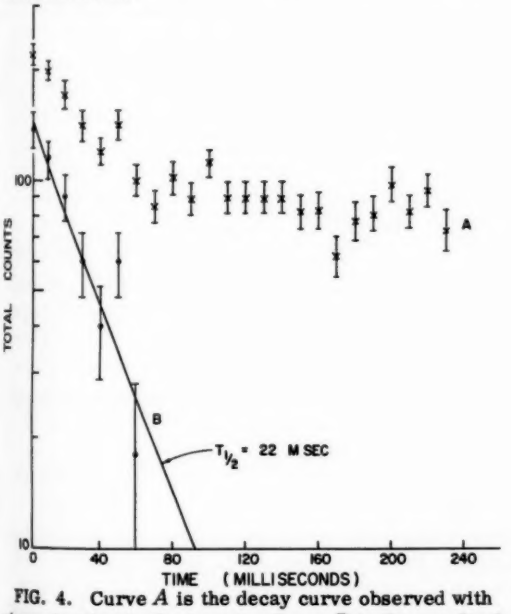

FIG. 4. Curve A is the decay curve observed with pions stopping in the target. Curve B is the result of subtracting the long-lived background. Each point represents the total number of counts obtained in a 10-msec interval.

total of 4×10^4 counts on the scalers which gave asymmetries $\epsilon_F = -(2.36\pm1.58)\%$, and $\epsilon_B = +(2.88\pm1.30)\%$. If we represent the polarized forward and backward intensities by $(1-a)$ and $(1+a)$, respectively, these results combined yield $a = +(2.68\pm1.00)\%$. Thus the B^{12} nuclei were polarized in the forward direction. These values do not take into account the dilution of the scaler asymmetries by the background. When correction is made for this effect, a is increased to $(4.3\pm1.7)\%$ for the B^{12}. Analysis of the time interval meter data gave agreement with this value, and indicated that the asymmetry was due solely to the B^{12} component. The value of a is consistent with the theoretical predictions,[1] and indicates positive helicity for the negative muons. This is in agreement with experiments on polarization of electrons and positrons from μ-decay,[6] and with the predictions of the law of conservation of leptons.

A measurement was also made of the asymmetry of β-rays from the activities produced by stopping pions using the same timing sequence as that for the B^{12} produced by muons. The result was $a_\pi = -(1.10\pm1.26)\%$.

Further measurements to improve the accuracy are in progress.

We wish to express our thanks to B. L. Bloch, E. J. Donelan, and E. J. R. Maier for assistance in many phases of the experiment, and to many members of the laboratory for helpful discussions and criticism. We also wish to thank R. A. Dannels for aiding in the data reduction and Dr. W. I. Goldburg for proton relaxation measurements.

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*This work was supported in part by the Atomic Energy Commission.

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[1] Jackson, Treiman, and Wyld, Phys. Rev. 107, 327 (1957).

[2] T. N. K. Godfrey, thesis, Princeton, 1954 (unpublished); H. V. Argo et al., J. Fetkovich et al., and W. A. Love et al., at the Gatlinburg Conference on Weak Interactions, October, 1958 [Argo, Harrison, Kruse, and McGuire, Phys. Rev. (to be published); J. Fetkovich et al., Bull. Am. Phys. Soc. Ser. II, 4, 81 (1959); W. A. Love et al., Bull. Am. Phys. Soc. Ser. II, 4, 81 (1959)].

[3] J. O. Burgman et al., Phys. Rev. Lett. 1, 469 (1958).

[4] F. Ajzenberg-Selove (private communication).

[5] See, for example, J. I. Friedman and V. L. Telegdi, Phys. Rev. 105, 1681 (1957).

[6] Culligan, Frank, and Holt, at the Gatlinburg Con-

ference on Weak Interactions, October, 1958 [Bulletin of the American Physical Society Ser. II, 4, 81 (1959)]; Macq, Crowe, and Haddock, Phys. Rev. 112, 2061 (1958).

SEARCH FOR AN 80-msec ACTIVITY IN LIGHT NUCLEI*

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In view of the report[1] of an 80-msec half-life activity produced by the capture of π^- mesons in carbon we have considered various possible products including He^7 and low-lying isomeric states, particularly[2] in Be^{11}. Such activities, if they exist, should also be formed by (n,p) reactions with fast neutrons.

Samples of 0.5-1.5 g of Li, Be, B, B^{10}, and C were exposed to 15.5-Mev neutrons resulting from a 20-μa beam of 600-kev deuterons incident on a thick Zr-tritium target. In the first set of experiments a pivoted arm moved the sample in 0.1 sec a distance of 2 feet from the neutron source to a position in front of a paraffin-shielded Pilot-B scintillator, 2 inches thick by 3 inches in diameter, and at the same time the beam was removed from the target. Backgrounds here were relatively high, so for the second series of tests we moved the detector to the shielded control room and the samples were transferred pneumatically in $\frac{1}{3}$ sec in thin-walled aluminum "rabbits." Beta rays of >4 Mev were accepted and their decay was displayed on a 100-channel analyzer adapted to serve as a time-channel analyzer. 21-msec B^{12} produced in the crystal was detected in the first series of experiments. Other activities observed include 0.17-sec Li^9 from the Be sample, and 0.8-sec Li^8 plus 13.6-sec Be^{11} from the B sample.

No other activities between ~0.05 and 10 sec were found. On the basis of the known 30-mb cross section for producing Li^8 by the $B^{11}(n,\alpha)Li^8$ reaction at $E_n = 15.5$ Mev, we find that the upper limits for cross sections of (n,p) reactions in the various nuclei, leading to the formation of a high-energy beta-ray emitter of 80-msec half-life, are (in millibarns) as follows: Li-1.2; Be-4; B^{11}-50; B^{10}-0.8; and C-2. Cross-section limits for producing nuclides of up to 10-sec half-life, other than those already known, are within an order of magnitude of these figures. The $Be^9(n,p)Li^9$ reaction, which has not been reported previously, has a cross section of ~0.6 mb at $E_n = 15.5$ Mev according to our rough preliminary measurements.

Meanwhile we have learned that the 80-msec activity reported by the CERN group is not confirmed by Love et al.[3] in a similar experiment.

We are indebted to M. Goldhaber for suggesting this problem.

*Work done under the auspices of the U. S. Atomic Energy Commission.

†On leave from Institut de Recherches Nucléaires, Strasbourg, France.

‡On leave from the Hebrew University, Jerusalem, Israel.

[1]Burgman, Fischer, Leontic, Lundby, Meunier, and Stroot, Phys. Rev. Lett. 1, 469 (1958).

[2]A 13.6-sec half-life has been assigned to the ground state of Be^{11} by M. J. Nurmia and R. W. Fink, Phys. Rev. Lett. 1, 23 (1958), and by D. E. Alburger and D. H. Wilkinson, Phil. Mag. 3, 1332 (1958); D. H. Wilkinson and D. E. Alburger, Phys. Rev. (to be published).

[3]R. T. Siegel (private communication); see Love, Marder, Nadelhaft, Siegel, and Taylor, preceding Letter [Phys. Rev. Lett. 2, 107 (1959)].

$\overline{K}^0$-K^- MASS EXCESS*

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Using the charge exchange of K^- in a hydrogen bubble chamber, we find that $\overline{K}^0$ is 3.7 ± 0.7 Mev heavier than K^-.

The incident K^- ranged in energy from 10 to 150 Mev. We have observed 44 charge-exchange events,

$$K^- + p \rightarrow \overline{K}^0 + n, \qquad (1)$$

in which the $\overline{K}^0$ decays into $\pi^+ + \pi^-$. In five of these events there was a recoil proton consistent with production by the neutron in Eq. (1).

For each event the $\overline{K}^0$ mass was adjusted to give a best fit to both production and decay vertices (taken simultaneously); however, the uncertainties can be best understood in terms of a

slightly oversimplified discussion in which the two vertices are fitted separately. The fact that the decay vertex is very exothermic makes the fit insensitive to the assumed $\overline{K}^0$ mass. This means on the one hand that this vertex (taken alone) yields little mass information, but on the other hand it gives a $\overline{K}^0$ momentum p_0 almost independent of uncertainty in its mass.

Cases without associated recoil protons: The production vertex is next fitted by using the calculated p_0 and the momentum p_- of the slow incoming K^- as determined by its measured curvature. This fit of the low-energy, endothermic production vertex is very sensitive to the $\overline{K}^0$-K^- mass excess; typically a 1-Mev change in mass excess requires about a 5-Mev/c change in p_-. The principal uncertainty is in p_-, caused by multiple Coulomb scattering. There is in addition a generally smaller error in p_0. The total uncertainty computed from these effects agrees very well with the external consistency of the various events. A least-squares fit is made of p_-(computed)-p_-(measured), with the mass excess taken as the variable; we find

$$M_{\overline{K}^0} - M_{K^-}\,(\text{no recoil}) = 4.7 \pm 1.3 \text{ Mev} \quad (2)$$

As a check on the reliability of the measured K^- momentum for a particle that is losing momentum rapidly, we have taken a sample of K^- mesons that come to rest in the chamber (as evidenced by the collinear $\Sigma + \pi$ produced from captures at rest) and measured the momentum of the tracks to about 6 cm (150 Mev/c) from the stopping end. The residual range provides a precise measure of the momentum at the point where the curvature measurement is stopped. Comparing the momentum obtained by curvature with that obtained by residual range, we obtained, for this entire sample, p(curvature) - p (range) = $+1.5 \pm 4.1$ Mev/c, showing that possible systematic effects in the curvature measurement are small. This 4.1-Mev/c uncertainty has been folded into the mass uncertainty quoted above.

Cases with recoils: We found five cases with proton recoils associated with the neutrons in reaction (1). The recoils were found by computing the neutron direction and momentum for each event and searching the appropriate volume of the chamber for a consistent recoil. The mean free path for an n-p collision of sufficiently large angle to make a visible recoil is approximately 100 cm, leading us to expect 4.6 recoils among those events in which the neutron is sufficiently energetic to produce a visible proton recoil. A typical bubble chamber picture contains in addition about eight recoils from background neutrons traversing the chamber. The probability is approximately 10% that in the entire volume searched there should be a recoil which accidentally satisfies coplanarity and conservation of transverse momentum within two standard deviations and lies within a $\overline{K}^0$ mass range of 494 ± 7 Mev.

For an individual event with recoil, determination of the mass difference becomes much more precise, because it does not require a measurement of momentum by curvature in a sensitive way, but only a measurement of proton range and angles. In Table I we have listed the five events in which a recoil was observed, along with the best-fit mass difference calculated for each event. The first event is clearly inconsistent with Eq. (2), and therefore we feel justified in identifying this as an accidental recoil.

The combined value of the no-recoil result and the four recoil events is 3.7 ± 0.7 Mev. If we take the K^- mass to be the same as the K^+, i.e., 494.0 ± 0.2 Mev,[1] the mass of the $\overline{K}^0$ is then 497.7 ± 0.8 Mev.[2] Assuming equality of the K^0 and $\overline{K}^0$ masses, one can combine this measurement with those listed in reference 2 to obtain a mass of 497.9 ± 0.6 Mev.

The fact that the $\overline{K}^0$ is heavier than the K^- is rather surprising. If they are members of a charge doublet, as commonly assumed, then one might expect for spinless particles that the charged member should be heavier, although no general proof of this statement is known.[3] On the other hand, if the charged K and neutral K are not members of a doublet--as, for example, in the theory of Pais[4]--then their masses need bear no relation to each other.

We wish to thank Professor Luis W. Alvarez

Table I. Best-fit mass differences for five events in which recoils were observed.

Event	$M_{\overline{K}^0}$-M_{K^-} (Mev)
1	-0.3 ± 1.2
2	$+4.1 \pm 1.3$
3	$+2.5 \pm 1.2$
4	$+9.0 \pm 7.0$
5	-5.0 ± 12.0
Average of 2-5	$+3.3 \pm 0.9$

for his advice and encouragement, and members of the bubble chamber and scanning staff for their assistance.

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*This work was done under the auspices of the U. S. Atomic Energy Commission.

[1]Cohen, Crowe, and DuMond, Nuovo cimento 5, 541 (1957).

[2]This may be compared to the following neutral K mass measurements (in Mev): 491.3±4 [Arnold, Martin, and Wyld, Phys. Rev. 100, 1545 (1955)]; 493.3±7.5 [Thompson, Burwell, and Huggett, Suppl. Nuovo cimento 4, 286, (1956)]; 501.9±5.3 [Fretter, Friesen, and Lagarrigue, Suppl. Nuovo cimento 4, 539 (1956)]; 500.8±7.7 [Fowler, Maenchen, Powell, Saphir, and Wright, Phys. Rev. 103, 208 (1956)]; 496.3±4 [D'Andlau, Armenteros, Astier, DeStaebler, Gregory, LePrince-Ringuet, Muller, Peyrou, and Tinlot, Nuovo cimento 6, 1135 (1957)]; 499.8±5.1 [Baxter H. Armstrong, University of California Radiation Laboratory Report UCRL-3470 (1956) (unpublished)]; 498.8±1.1 [Crawford, Cresti, Good, Stevenson, and Ticho, following Letter, Phys. Rev. Lett. 2, 112 (1959)].

[3]It has been observed by S. Gasiorowicz and A. Petermann [Phys. Rev. Lett. 1, 457 (1958)] that, in perturbation theory, electromagnetic mass corrections for spinless particles should lead to the following relation for pions and K mesons: $(M_{K^+}-M_{K^0})/(M_{\pi^+}-M_{\pi^0}) = M_\pi/M_K$, contrary to the results presented here.

[4]A. Pais, Phys. Rev. 112, 624 (1958).

K^0-K^+ MASS EXCESS

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In a study of associated production by 1.12-Bev/c negative pions incident on a liquid hydrogen bubble chamber, we find a discrepancy between the K^0 and K^+ production and decay dynamics, if we use currently accepted values for the masses involved. The discrepancy is most easily resolved if the mass of the K^0 exceeds that of the K^+ by about 5 Mev.

The experiment is sensitive to the mass values, in spite of the high energies involved, because one is working close to associated production threshold. The method follows.

The average beam momentum $\bar{P}_{\text{inc}}$ and its rms half-width $\delta\bar{P}_{\text{inc}}$ are determined from 16 selected events of the type

$$\pi^- + p \rightarrow \Sigma^- + K^+, \tag{1}$$

in which the K^+ stops in the chamber, and thus has its momentum accurately determined by range. Using currently accepted mass values (Table I), we obtain the result

$$\bar{P}_{\text{inc}} = 1123.5 \pm 2.4 \text{ Mev}/c,$$
$$\delta\bar{P}_{\text{inc}} = 8.2 \pm 2 \text{ Mev}/c, \tag{2}$$

where the error in beam momentum includes propagation of the (relatively small) contributions from the uncertainties in the π^-, Σ^-, and K^+ masses, as given in Table I.

For each event of the type

$$\pi^- + p \rightarrow \Lambda + K^0, \tag{3}$$

with one or both of the Λ and K^0 decaying in the chamber, the momentum of the incident pion is calculated from the production and decay dynamics. This is done as a routine part of the data analysis. When 494 Mev is used as the K^0 mass, the beam momenta thus obtained are systematically about 1% too low, as compared with the Σ^-K^+ result in Eq. (2). This is the discrepancy referred to at the beginning.

The data were accordingly rerun through an improved IBM program which includes the added feature that it propagates errors in the beam momentum. In this program a K^0 mass of 498.0 Mev was used; this was close to the value suggested by the discrepancy. The discrepancy was thereby of course very much reduced.

From the events in which both the Λ and K^0 decay in the chamber we then select those with small errors, $\delta\bar{P}_{\text{inc}} < 20$ Mev/c, on the incident pion momentum. It is a simple matter to transform the residual discrepancy between the resulting $\bar{P}_{\text{inc}}$ and the value obtained from Eq. (2) into a deviation of M_{K^0} from 498, and propagate the errors thereon.

There were 34 ΛK^0 events that passed the selection criterion. From these events, and using the value $M_{K^0} = 498$ in the IBM program, we obtain $\bar{P}_{\text{inc}} = 1121.1 \pm 2.3$. In order to agree with the Σ^-K^+ result (2), the K^0 mass must be increased by an additional 0.8 Mev. We finally obtain the result

$$M_{K^0} - M_{K^+} = 4.8 \pm 1.1 \text{ Mev}, \tag{4}$$

Table I. Mass values and their derivatives used in computing K^0-K^+ mass difference.[a]

	π^-	Σ^-	K^+	Λ
Mass, M	139.63	1196.50	494.00	1115.20
δM_{rms}	0.06	0.50	0.20	0.14
$\partial M_{K^0}/\partial M$	-0.03	0.73	1.06	-0.54

[a]The first two lines in Table I contain mass values and their rms errors, in Mev, used in our determination of K^0-K^+ mass excess. They correspond to Table I of M. Gell-Mann and A. H. Rosenfeld, Annual Review of Nuclear Science (Annual Reviews, Inc., Stanford, 1957), Vol. 7. The last line gives partial derivatives of our K^0 mass value with respect to the other assumed masses, in Mev per Mev. The smallness of the dependence, -0.03, on the π^- mass reflects the fact that the π^- mass "cancels" to first order in the K^0-K^+ mass difference. The nearly 1:1 (1.06) dependence on the K^+ mass reflects the fact that to first order our K^0-K^+ mass excess is independent of the value assigned to the K^+ mass. The contribution to the error in M_{K^0}-M_{K^+} that arises from errors in the masses listed in this table is obtained by multiplying the second and third rows of the table together, with the usual squaring and adding of terms.

where the error includes the uncertainties in $\bar{P}_{\text{inc}}$ as determined from the 34 ΛK^0 events, and from the 16 ΣK^+ events, and also includes the errors in the π^-, Σ^-, K^+, and Λ masses. We emphasize that it is the mass difference which is determined--the result in Eq. (4) would be unchanged to first order even if the currently accepted K^+ mass were found to be too low.

We next turn to the question of possible systematic errors. The determination of beam momentum for a ΛK^0 event depends on measurements of the momentum and laboratory-system production angle of the K^0. (The Λ production and decay are used to check the event for consistency, but are not used in the actual determination of P_{inc}.) The K^0 momentum in turn is determined almost completely by angle measurements on the decay pions, and is quite insensitive to the measurement of their curvature in the 11-kilogauss magnetic field. (In other words, the magnitude and error of the K^0-K^+ mass difference obtained above would have been essentially unchanged if the magnetic field had been turned off.) The question then becomes one of systematic errors in angle measurements.

The best confirmation that angle measurements are not significantly wrong comes from the work of Baggett and McCormick.[1] In order to determine the incident pion momentum, they analyzed elastic π -p scattering events in which the proton stops in the chamber. Their events are taken from the same series of bubble chamber picture that contain the associated production events. The beam momentum is essentially given by proton range and pion angle. Their result, $\bar{P}_{\text{inc}} = 1135 \pm 10$ Mev/c, agrees (within the errors) with our $\Sigma^- K^+$ result. This agreement means that angles are measured very well indeed; one standard deviation on their result represents a systematic shift of 0.1 degree in pion angle.

The angle errors needed to reconcile our data with $M_{K^0} - M_{K^+} = 0$ are, however, about 1 degree. If we weigh in the possibility of small systematic angle errors of the order of 0.1 degree, the rms error in Eq. (4) is increased from 1.1 to 1.14 Mev.

The masses used in this determination, and the partial derivatives of M_{K^0} with respect to them, are given in Table I.

For a compilation and weighted average of our K^0 mass and K^0 and $\bar{K}^0$ mass values obtained by other workers, we refer to the accompanying Letter by Rosenfeld, Tripp, and Solmitz.[2]

Finally we remark that our K^0-K^+ mass excess has the opposite sign, and roughly three times the magnitude expected[3] if (a) the K^0 and K^+ are members of the same multiplet, (b) the mass difference is purely electromagnetic, and (c) the strong interactions are assumed to have a negligible effect. The suggestion of Pais, that the K^0 and K^+ may have opposite intrinsic parity,[4] would naturally invalidate assumption (a).

We would like to thank Roger L. Douglas and George R. Kalbfleisch for their assistance in analysing data, John Dardis and Paul Kenny for interesting discussions, and Luis W. Alvarez for his advice and encouragement.

[1]Private communications from Lee Baggett, Jr., and Bruce H. McCormick, and University of California Radiation Laboratory Report UCRL-8302 (unpublished).

[2]A. H. Rosenfeld, R. D. Tripp, and F. Solmitz, preceding Letter [Phys. Rev. Lett. 2, 110 (1959)].

[3]S. Gasiorowicz and A. Petermann, Phys. Rev. Lett. 1, 457 (1958).

[4]A. Pais, Phys. Rev. 112, 624 (1958).

EXPERIMENTAL DETERMINATION OF THE Λ SPIN*

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Lee and Yang have recently proposed a method of determining the Λ spin.[1] In this Letter we apply their method to a sample of 614 decays of the type

$$\Lambda \rightarrow \pi^- + p. \quad (1)$$

The Λ's were produced in our 10-inch liquid hydrogen bubble chamber via the reaction

$$\pi^- + p \rightarrow K^0 + \Lambda, \quad (2)$$

by incident pions of momenta 1.23 Bev/c (184 events), 1.12 Bev/c (253 events), 1.09 Bev/c (53 events), 1.03 Bev/c (94 events), and 0.95 Bev/c (30 events).

The beauty of the Lee-Yang method is that it makes possible an absolute determination of the Λ spin without assumptions other than that angular momentum is conserved in the Λ decay.[2] This can be contrasted with the method of Adair,[3] in which details of the production reaction (2) must be considered in order to reach a conclusion about the Λ spin. In particular, assumptions must be made as to the final orbital-angular-momentum states present in the $K^0+\Lambda$ system. Similarly, assumptions must be made about the K^0 spin, and a selection of the data made accordingly; the lower the assumed K^0 spin, the more data one can use. For instance, Eisler et al.[4] have applied the Adair analysis to their associated production data and have concluded that the Λ spin is $\frac{1}{2}$, provided that the K^0 spin is zero, and provided that only S, P, and D waves are important in the $K^0+\Lambda$ system.

The disadvantage of the Lee-Yang method is that a very large amount of data is needed in order to achieve conclusive results.[5]

The Lee-Yang method depends for its success upon the very large up-down asymmetry found[6-8] in the parity-nonconserving decay (1) of Λ's produced in reaction (2). In principle, one starts with any collection of Λ's, chooses a quantization direction in any way that is independent of the decay (1), and then examines the decay distribution $W(\xi)d\xi$, where ξ denotes the cosine of the angle between the (negative) decay pion and the quantization direction—all in the Λ rest frame. For a Λ spin J, $W(\xi)$ is in general a polynomial with powers up to and including ξ^{2J}. One might suppose that a given experimental ξ distribution could always be more easily fitted to large spin J than to small spin, since then more coefficients are available for "curve fitting." However, angular-momentum conservation severely constrains the coefficients. In fact, a very large up-down decay asymmetry cannot be achieved with a large spin. (This can be understood qualitatively through the observation that, classically, the disintegration of a system of high spin tends to yield fragments moving in the equatorial plane rather than towards the poles.)

The constraints on the shape of $W(\xi)$ can be summarized through the Lee-Yang test functions $T_{J,M}(\xi)$, all of which satisfy the inequality

$$\langle T_{J,M} \rangle \leq 1, \quad (3)$$

where the bracket denotes averaging over the decay distribution $W(\xi)$, and where $M = J, J-1, \ldots, -J$.

For spin $J=\frac{1}{2}$, the decay distribution is

$$W(\xi)d\xi = \tfrac{1}{2}\, d\xi(1+a\xi). \quad (4)$$

The test functions are $T_{1/2,\pm 1/2} = \pm 3\xi$, so that one has $\langle T_{1/2,\pm 1/2} \rangle = \pm a$, and the Lee-Yang inequality (3) reduces to $-1 \leq a \leq 1$.

For Λ spin $J=\frac{3}{2}$, the decay distribution is a cubic in ξ. The four test functions are

$$T_{3/2,3/2} = 9P_1(\xi) + 5P_2(\xi) - (7/3)P_3(\xi), \quad (5)$$

$$T_{3/2,1/2} = 3P_1(\xi) - 5P_2(\xi) + 7P_3(\xi), \quad (6)$$

and $T_{3/2,-3/2}$ and $T_{3/2,-1/2}$, which are obtained by substituting $-\xi$ for ξ in Eqs. (5) and (6). The $P_k(\xi)$ are Legendre polynomials. Similar test functions are constructed for spin $\frac{5}{2}$, $\frac{7}{2}$, etc.

To illustrate the method, suppose the Λ spin were really $\frac{1}{2}$, and that we had a decay sample with the maximum possible asymmetry, that is with $a=1$. Then we would find (with enough data so that statistical fluctuations were negligible) $\langle T_{3/2,\pm 3/2} \rangle = \pm 3a = \pm 3$, and $\langle T_{3/2,\pm 1/2} \rangle = \pm a = \pm 1$. The first of these fails to satisfy the inequality (3), and spin $\frac{3}{2}$ would thereby be ruled out. The other three spin-$\frac{3}{2}$ test functions satisfy the inequality and yield no information. We notice that a sample of spin-$\frac{1}{2}$ Λ's having $|a| \leq \frac{1}{3}$ would satisfy all four of the spin-$\frac{3}{2}$ Lee-Yang inequalities, and therefore would be useless for application of the method.

We now consider the way in which the sample of Λ decays is obtained. For the method to suc-

ceed, the preceding example shows that one needs a sample with a large asymmetry with respect to the quantization axis. On the other hand, in order to obtain an unbiased sample it is of primary importance not to select the Λ's or the quantization axis in a manner which involves "peeking" at the decay ξ values. With these considerations in mind, we naturally choose a priori a quantization axis perpendicular to the production plane of reaction (2), along $\vec{P}(\pi \text{ incident}) \times \vec{P}(\Lambda)$. We naturally exclude, a priori, from our sample any Λ's produced via $\pi^- + p \to K^0 + \Sigma^0$, $\Sigma^0 \to \gamma + \Lambda$. The question arises whether we should include the entire range of c.m. angles in the production reaction (2), or, as suggested by Lee and Yang,[1] include only a region centered at 90° (c.m.), where the polarization might be expected to be largest. We believe that it is very difficult to justify such a limitation, since in order to decide on the range of angles to be included one becomes involved either in a posteriori "peeking at the data," or in making implicit a priori assumptions as to the maximum angular complexity and thus as to the maximum number of partial waves involved in the $K^0 + \Lambda$ state. The former biases the distribution, and the latter spoils the beauty of the assumption-free Lee-Yang method. We therefore include the entire range of production angles. Similarly, the question arises whether we should include all incident-pion production energies, or only those in which the decay asymmetry appears to be largest. Since we have no a priori knowledge as to the energy dependence of the Λ polarization, we would be at the mercy of statistical fluctuations, with a consequent large chance for bias, if we excluded some datum because of its small observed up-down asymmetry. We therefore include, a priori, all production energies.

Finally we present the results. For each event, ξ is obtained from detailed dynamical analysis. Then $T_{3/2,3/2}(\xi)$ is calculated for each event, by using Eq. (5). Then (omitting the $\frac{3}{2}, \frac{3}{2}$ subscripts), $\langle T \rangle = (1/N)\sum T(\xi) \pm [(1/N)(\langle T^2 \rangle - \langle T \rangle^2)]^{1/2}$, where $\langle T^2 \rangle = (1/N)\sum T^2(\xi)$, and $N = 614$ is the total number of events. The other test functions are calculated analogously. Our 614 events yield

$$\langle T_{1/2,1/2} \rangle = 0.57 \pm 0.066,$$

$$\langle T_{3/2,3/2} \rangle = 1.77 \pm 0.244,$$

$$\langle T_{5/2,3/2} \rangle = 2.99 \pm 0.408.$$

Thus a Λ spin of $\frac{1}{2}$ easily satisfies the Lee-Yang inequality $\langle T_{J,M} \rangle \leq 1$, while spin $\frac{3}{2}$ fails to satisfy it by $(1.77-1)/(0.244) = 3.16$ standard deviations, and spin $\frac{5}{2}$ fails by 4.88 standard deviations.[9]

In addition to satisfying the spin-$\frac{1}{2}$ Lee-Yang inequality, $3\langle \xi \rangle \leq 1$, the ξ distribution must be linear [Eq. (4)], for spin $\frac{1}{2}$. Figure 1 shows a histogram of the experimental distribution. The straight line is a least-squares best fit and corresponds to the slope $a = 0.57$. Application of the χ^2 test to the fit yields $\chi^2 = 6.70$, with an "expected" value of $10 - 2 = 8$. This corresponds to a χ^2 probability of 57% for a fit this bad or worse. The data thus fit a linear distribution (4) very well indeed.

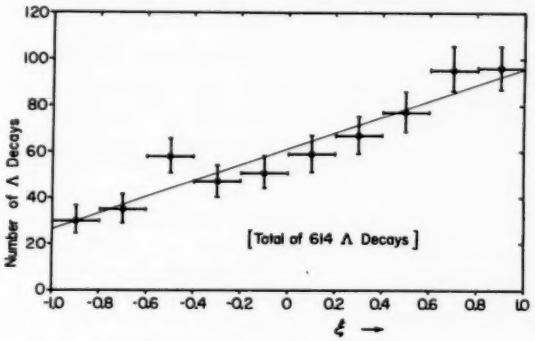

FIG. 1. "Up-down" distribution for $\Lambda \to \pi^- + p$.

Lastly, we have performed a control experiment, in order to search for possible hidden systematic errors in our determination of the Λ decay ξ distribution. Namely, we have determined, in exactly the same way as for the Λ's, the "ξ distribution" of a sample of 547 K^0 decays of the type $K^0 \to \pi^+ + \pi^-$. This sample includes substantially all our K^0 charged decays, from all production energies and from K^0's produced in association with Σ^0's as well as with Λ's. Figure 2 shows the results. As expected if there are

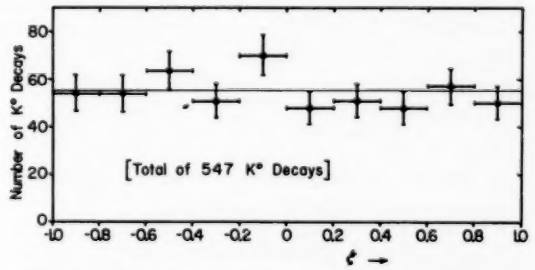

FIG. 2. Control experiment: up-down decay distribution for $K^0 \to \pi^- + \pi^+$.

indeed no biases, the ξ distribution is flat.[10]

We finally conclude, free from assumptions, that the Λ spin is $\frac{1}{2}$.

We are indebted to Frank T. Solmitz for many fruitful discussions, to George R. Kalbfleisch and Roger L. Douglas for their help in analyzing much of the data, and to Luis W. Alvarez for his continued guidance and support.

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[1]T. D. Lee and C. N. Yang, Phys. Rev. 109, 1755 (1958). An extension to include the polarization of the decay proton is discussed by Durand, Landovitz, and Leitner, Phys. Rev. 112, 273 (1958).

[2]The decay proton and pion spins are assumed to be 1/2 and zero. For a Λ spin of 1/2, the final state in decay (1) is an interfering mixture of states $S_{1/2}$ and $P_{1/2}$, for spin 3/2 a mixture of $P_{3/2}$ and $D_{3/2}$, etc.

[3]R. K. Adair, Phys. Rev. 100, 1540 (1955).

[4]F. Eisler et al., Nuovo cimento 7, 222 (1958).

[5]In the 1958 Annual International Conference on High-Energy Physics at CERN, edited by B. Ferretti (CERN, Geneva, 1958), we reported a Lee-Yang type analysis of 237 of our events at 1.12 Bev/c, plus a total of 162 events obtained by private communications from M. Schwartz and D. Glaser. The result was that the evidence against spin 3/2 was not conclusive.

[6]Crawford, Cresti, Good, Gottstein, Lyman, Solmitz, Stevenson, and Ticho, Phys. Rev. 108, 1102 (1957).

[7]Eisler, Plano, Prodell, Samios, Schwartz, Steinberger, Bassi, Borelli, Puppi, Tanaka, Woloschek, Zoboli, Conversi, Franzini, Mannelli, Santangelo, Silvestrini, Brown, Glaser, Graves, and Perl, Phys. Rev. 108, 1353 (1957).

[8]L. B. Leipuner and R. K. Adair, Phys. Rev. 109, 1358 (1958).

[9]Higher spins that 5/2 fail by an even larger margin. Also, as was expected from the illustrative example, the other spin-3/2 test functions do satisfy the Lee-Yang inequality, and therefore yield no information.

[10]Application of the χ^2 test to the horizontal straight line drawn through the experimental points yields $\chi^2 = 7.66$, where $10-1=9$ is "expected." This corresponds to a χ^2 probability of 56%, and is thus a very good fit.

PHOTOGRAPHY OF COSMIC RAYS IN A LUMINESCENT CHAMBER*

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There is an increasing need in high-energy physics for a counter-controlled, high time-resolution device which records the tracks of particles. Such a device is the luminescent chamber,[1-4] in which the light from a charged particle passing through a scintillator is amplified in an image preserving manner and then photographically recorded. Chambers of useful size can be made using either a single scintillator crystal[1-4] or a bundle of scintillator filaments.[3] The major problem of obtaining sufficiently high light amplification had until now only been solved by Zavoisky and co-workers.[1,2] Unfortunately very little technical detail on this work has been available.

The authors have succeeded in consistently photographing minimum-ionizing cosmic rays in a sodium iodide crystal using available electronic and optical components. We believe that this is the first time that either cosmic rays or minimum-ionizing tracks have been recorded in a luminescent chamber. With modifications this apparatus can be used for luminescent chamber experiments in high-energy particle physics.

The system, Fig. 1, consists of three image-intensifier tubes, optically coupled to each other and to the crystal and film with refractive lenses. The total amplification of the system corresponds to about $2-8\times10^4$ photons falling on the film for each photon which strikes the first photocathode of the system. The first and third tubes are on continuously, and the second tube is pulsed on for one millisecond when a photomultiplier detects a cosmic ray passing through the crystal. The crystal is 3.8 cm in diameter and 1.9 cm thick. The lens viewing the crystal collects 0.5% of the light emitted by the particle.

Figure 2 shows a cosmic-ray track and also some image-tube background noise. Our optics would focus 530 quanta from each centimeter of minimum-ionizing track in the crystal onto the first image-tube cathode, releasing about 50 photoelectrons. This yields 150 photoelectrons from the entire track. In the photograph, the track is 3 cm long and 1 mm thick. This thick-

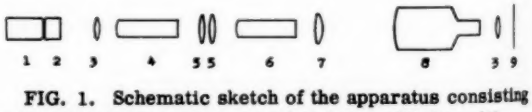

FIG. 1. Schematic sketch of the apparatus consisting of the following components: (1) photomultiplier; (2) 3.8-cm diameter by 1.9-cm thick NaI(Tl) crystal; (3) Erfle eyepiece lenses; (4) RCA C 73458A two-stage image tube; (5) American Optical eyepiece lenses; (6) RCA C 73458 two-stage image tube; (7) Kodak 110 mm, $f/0.75$ lens; (8) Westinghouse WX3897 one-stage image tube; (9) Kodak Royal X-Pan film.

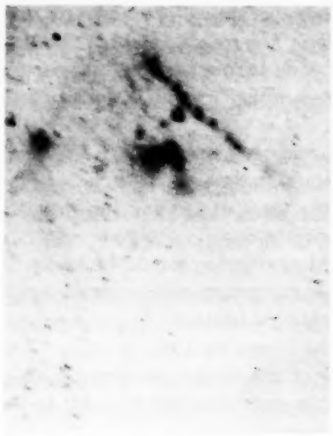

FIG. 2. Track photograph. The track is about 3 cm long in the crystal and is due to a minimum-ionizing cosmic ray. The three black dots, one near the track and two at the center, are due to clumps of electrons ejected from the first cathode by ion bombardment.

ness, corresponding to a 1 mm over-all resolution, is a result of the depth of field in the crystal as well as the combined resolutions of the image tubes and optics. Thus on the average, each resolution element corresponds to five photoelectrons; and the statistical fluctuations result in an inhomogeneous track. This inhomogeneity appears greater in Fig. 2 because of the contrast enhancement occurring in the photographic copying techniques. The resolution of the system is too poor and the gain too low to permit observation of individual photons at present. The background in Fig. 2 is due to positive-ion bombardment of the first cathode and appears whether or not the tube gate signal is derived from a cosmic ray.

In 150 photographs taken from trigger pulses corresponding to greater than 3 Mev energy loss in the crystal, 54 clearly identified tracks were seen. Other tracks were present in about 50 pictures but were questionable because they were too short (corresponding to an angle through the crystal close to the lens axis) or were obscured by the image-tube background. From the known spectrum of cosmic-ray μ mesons[5] the majority of all of these tracks must be minimum-ionizing μ mesons. A few might be high-energy and consequently minimum-ionizing electrons. The tracks can also be observed visually by looking at the last anode of the system.

*This work is being supported by the Office of Naval Research, by the Michigan Memorial Phoenix Project of the University of Michigan, and by Project Michigan for the U. S. Army Signal Corps, under a Department of the Army contract, administered through the University of Michigan Research Institute.

[1]Zavoisky, Smolkin, Plakhov, and Butslov, Doklady Akad. Nauk S.S.S.R. 100, 241 (1955).

[2]Zavoisky, Butslov, Plakhov, and Smolkin, Atomnaya Energ. 1, 34 (1956) [translation: J. Nuclear Energy 4, 340 (1957).

[3]G. T. Reynolds and P. E. Condon, Rev. Sci. Instr. 28, 1098 (1957).

[4]L. W. Jones and M. L. Perl, Rev. Sci. Instr. 29, 441 (1958).

[5]J. G. Wilson, Nature 158, 414 (1946).

K-MESON–NUCLEON TOTAL CROSS SECTIONS FROM 0.6 TO 2.0 Bev*

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Experiments by various groups[1] have shown the K^+-proton total cross section to be about 15 mb in the 100-200 Mev range and suggest an increasing cross section up to 300 Mev, the highest energy previously measured. We have explored the energy region beyond this with the hope that the total cross section might exhibit a resonance, analogous to the pion-nucleon behavior, or show some structure around the threshold for, e.g., 2π-meson production.

A differential Čerenkov counter of new design was used to select K mesons from a momentum-analyzed beam by their characteristic velocity. The counter will be described in detail elsewhere; the important features are as follows:

(a) High selectivity—at a typical adjustment the counting efficiency is 0.7 for collimated particles of $\beta = 0.98$. The efficiency drops to half this for a change $\Delta\beta = \pm 0.003$ and drops to about 0.001 for $\Delta\beta = \pm 0.015$.

(b) Variable operating range—the Čerenkov medium, fluorochemical FC75,[2] was chosen for its low critical pressure. By suitable choice of pressure and temperature, the index of refrac-

tion can be varied continuously from 1.00 to 1.28, the index of the liquid at room temperature. Thus the counter can be tuned continuously over the range from $\beta = 1$ to $\beta = 0.79$.

Figure 1 shows the beam geometry and counter array which was used at the Bevatron. A K-meson beam was defined in space by a triple coincidence C_1C_2S, momentum-selected by magnet M_1, and velocity-selected by counter C_1; the momentum and velocity selections were repeated by M_2 and C_2 to insure a high rejection for π mesons. The scintillator S was required to reduce the accidental coincidence rate. Those particles which were not scattered in the absorber were detected by the scintillator D.

The experiment consisted of measuring the attenuation in good geometry by a 5-ft long liquid hydrogen target.[3] The attenuation came directly from the ratio of threefold (C_1C_2S) to fourfold (C_1C_2SD) coincidences and was normalized to the attenuation of a dummy empty target. The scintillator M was put in a coincidence arrangement such that the efficiency of counter D was monitored continuously during the run. The K^+-neutron cross section was measured in the same geometry by the D_2O-H_2O difference. For 10^{10} protons on the target we obtained a yield at 1.75 Bev/c of about 6 K^+ mesons through the system, which had a momentum acceptance of $\pm 2\%$.

Figure 2 displays our data. No corrections have been made to the data and the errors shown are statistical only. The following considerations indicate that the required corrections are small:

(a) The separation of K^+-mesons from other particles was nearly complete. The background was 3% or less throughout the range, except for the 0.58-Bev point, at which the background was about 10%.

(b) The geometry was good in that the defining counter D subtended 0.018 sterad at the scatterer. If the total cross section were all elastic and isotropic in the c.m. system the correction for forward scattering would be about 1%. From a rough measurement of the small-angle scattering it appears unlikely that the correction should be as large as 10%.

(c) Corrections to the neutron cross section for shielding and coherent effects[4] in the deuteron

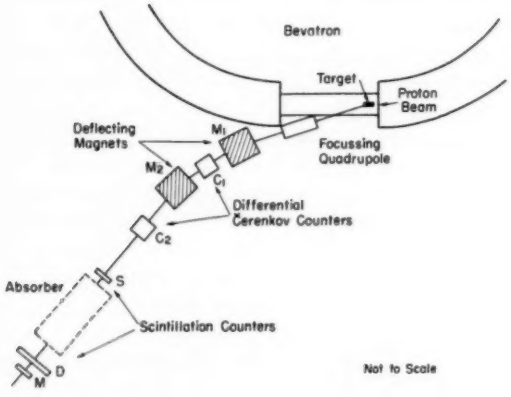

FIG. 1. Experimental arrangement. All counters are 4-in. diameter except D which is $8\frac{1}{4}$-in. The liquid hydrogen absorber is 8-in. diameter by 62-in. long and is centered 65 in. from counter D. Total flight path from Bevatron target to counter D is 17.4 meters. Shielding is not shown.

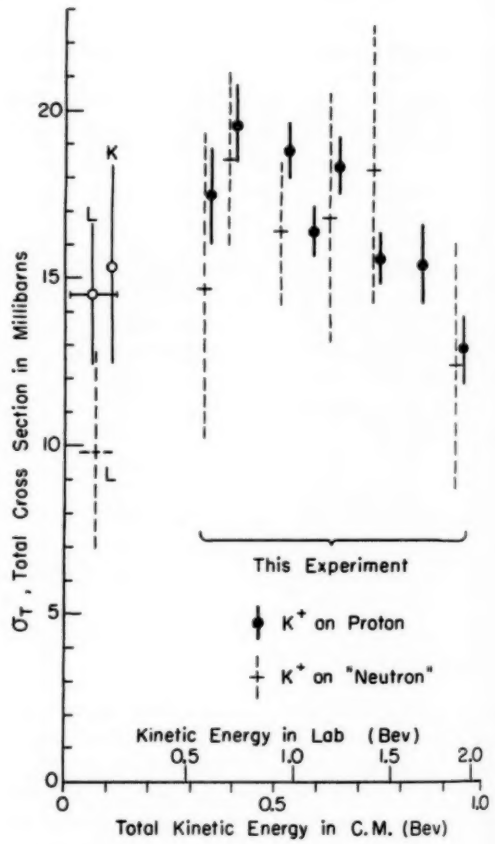

FIG. 2. Total cross sections for K^+ mesons on protons and "neutrons." The "neutron" points are the measured deuteron-proton difference. The results of Lannutti (L) and Kerth (K) are shown as representative of the low-energy data.[1]

would be about one millibarn if the scattering were largely inelastic or elastic and isotropic. The effect otherwise could be much larger depending on the real parts of the neutron and proton forward scattering amplitudes, which are unknown.

From the attenuation of the H_2O target we obtain for the total K^+-oxygen cross section (for absorption and diffraction scattering) at 0.69 Bev a value of 270 ± 20 mb which is 90% of geometric.

The K^--proton total cross section was measured at 0.96 and 1.23 Bev with the result 36 ± 5 and 44 ± 5 mb, respectively. From the decay in flight we obtained the lifetime $\tau_{K^+} = (1.21 \pm 0.06) \times 10^{-8}$ sec in good agreement with previous measurements on stopped K's. By this method we measured the ratio of K^- to K^+ lifetime directly and got $\tau_{K^-}/\tau_{K^+} = 0.97 \pm 0.09$.

We note the following:

(a) The K^+-proton total cross section reaches a maximum of $4\pi(\hbar/m_K c)^2$ or about $\pi(\hbar/2m_\pi c)^2$ at $P_{\text{c.m.}}^* \approx m_K c$; i.e., a total c.m. kinetic energy of $m_K c^2$.

(b) The K^+-proton and K^+-neutron total cross sections are about the same in this energy range; the K^--proton cross section is roughly twice as large.

(c) The contribution to the dispersion relation as evaluated by Matthews and Salam[5] is not negligible in the total energy range above $2m_K$ as they assumed tentatively. Assuming from our results that $\sigma(K^-) - \sigma(K^+)$ in the energy range 2 to 5 m_K is 25 mb, we find a contribution of -0.5 in their units.

The results of a new measurement of the π^--proton and π^+- proton total cross sections and acknowledgments are given in an accompanying Letter.

* This work is supported in part by funds provided by the U. S. Atomic Energy Commission, the Office of Naval Research, and the Air Force Office of Scientific Research.

† On the staff of the University of California Radiation Laboratory during the performance of this experiment.

[1] The Proceedings of the 1958 CERN Conference summarizes the existing data with references to the original papers. See also J. E. Lannutti et al., Phys. Rev. 109, 2121 (1958).

[2] Produced by the Minnesota Mining and Manufacturing Company.

[3] Cork, Lambertson, Piccioni, and Wenzel, Phys. Rev. 107, 248 (1957).

[4] Glauber, Phys. Rev. 100, 242 (1955).

[5] P. T. Matthews and A. Salam, Phys. Rev. 110, 569 (1958); see also C. Goebel, Phys. Rev. 110, 572 (1958); K. Igi, Progr. Theoret. Phys. Japan 19, 238 (1958); M. L. Goldberger, 1958 Annual International Conference on High-Energy Physics at CERN, edited by B. Ferretti (CERN, Geneva, 1958;) S. F. Tuan, University of California Radiation Laboratory Report UCRL-8461, 1958 (unpublished).

PION-PROTON RESONANCE SCATTERING NEAR 900 Mev*

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The scattering of negative and positive pions by hydrogen in the laboratory energy range 470 to 1200 Mev has been investigated in order to clarify the broad resonance occurring near 900 Mev.[1] The experimental setups were those used also for K-meson scattering experiments.[2] The energy calibration in the first arrangement (dots in Fig. 1) was done by using a current-carrying wire through the entire system to establish magnet currents at 870-Mev lab pion energy. Other points were determined by using the empirical magnetization curves for quadrupole and deflecting magnets, the latter operating at about 8000 gauss. The second arrangement (crosses in Fig. 1) was calibrated by wire for 2-Bev pion energy. The overlap shows consistency in the energy calibration.

The target was five feet of liquid hydrogen. The final counter defined cones of half-angles 6° and 2.3° about the downstream and upstream ends of the liquid hydrogen, respectively. The focussing Čerenkov counters discriminated against muons and electrons in the beam. At 600-Mev lab pion energy the ratio of efficiency for muons to that for pions was about 5%; at 900 Mev, 20%. Since the total muon flux is not over about 10% of pion flux, the muon flux counted is 2% or less of pion flux. The efficiency for counting protons in the π^+ beam was negligible.

On the basis of pion photoproduction measurements[3,4] it had been suggested by Wilson[5] that there exists an additional resonance in the iso-

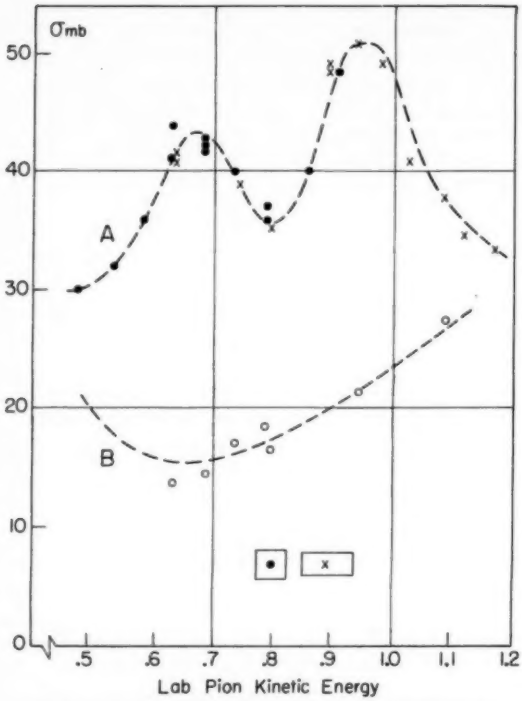

FIG. 1. Total cross sections for π^--p scattering (A) and for π^+-p scattering (B). Typical energy intervals are shown. (Energies in Bev.)

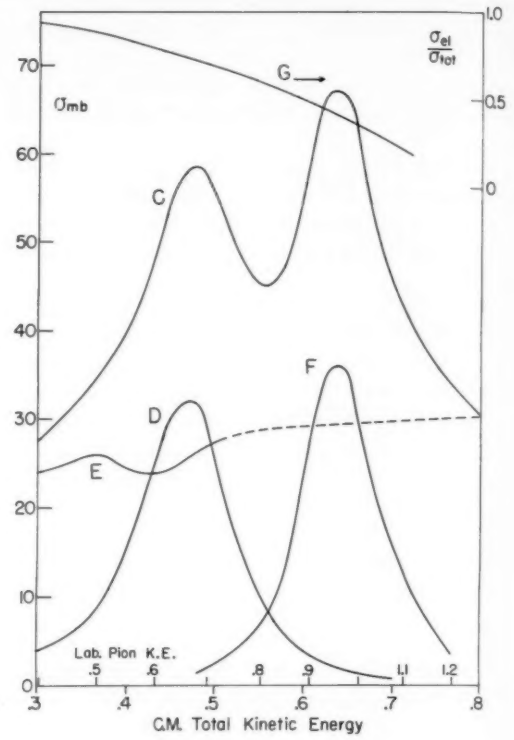

FIG. 2. Isotopic spin 1/2 total scattering cross section (C). The other curves are referred to in the text. The c.m. total kinetic energy is defined as the total energy in the c.m. system minus the rest energies of the proton and pion. (Energies in Bev.)

topic spin 1/2 state of the pion-proton system at a center-of-mass energy of about 600 Mev (pion laboratory kinetic energy 650 Mev) and having angular momentum 3/2.

In the π^- scattering [Fig. 1 (A)] our results show two clearly separated peaks at total center-of-mass energies of 615 ± 40 and 775 ± 40 Mev. There are no peaks in the positive-pion scattering in this region [Fig. 1 (B)], indicating these resonances are purely isotopic spin 1/2. The total isotopic spin 1/2 cross section is shown in Fig. 2 (C), in which the peak shapes have been adjusted for the experimental energy resolution indicated in Fig. 1.

The interpretation in terms of resonances in states of one total angular momentum depends on what assumption is made for the ratio σ_{el}/σ_{tot} for the resonant part of the intensity. The ratio σ_{el}/σ_{tot} for the sum of resonant and nonresonant intensity has been measured[6] to be 0.42 ± 0.02 at 950-Mev lab pion kinetic energy. If it is assumed that $\sigma_{inelastic}/\sigma_{total} = 1 - \sigma_{el}/\sigma_{tot}$ depends on the square of the excess kinetic energy in the center-of-mass system between the pion production threshold ($\sigma_{inelastic}/\sigma_{tot} = 0$) and the value 0.58 ± 0.02 at 950-Mev lab pion kinetic energy (G), and furthermore that this ratio holds for the resonant part, then the lower peak is consistent with a $j = 3/2$ resonance (D) calculated for a total width of 60 Mev. This interpretation leaves about 25 mb to be accounted for by an essentially energy-independent nonresonant interaction (E). An additional resonance in $j = 1/2$ at this energy cannot be excluded by present results.

The sum (615 ± 40 Mev) of the pion rest energy and the center-of-mass total kinetic energy at which the lower peak occurs is about 75 Mev higher than the corresponding energy at which the peak is observed in pion photoproduction according to the Calif. Inst. of Technology results.[4]

Subtracting a similar nonresonant intensity (E) at the upper peak leads to a peak value (F) of

about 36 mb. This is about twice that attributable to a $j=5/2$ resonance; however, the uncertainty in the magnitude of the nonresonant intensity, the possibility that σ_{el}/σ_{tot} is less than 0.4 for the resonant interaction, and the possibility of interference between the resonant and nonresonant amplitudes do not permit exclusion of $j=5/2$ for the upper resonance.

We are greatly indebted to Dr. E. J. Lofgren and the operating and engineering staff of the Bevatron for use of the Bevatron facilities, and to Dr. Bruce Cork for the help with the hydrogen target. The Sloan Foundation kindly supported one of us (D. H. Frisch) in the purchase of some equipment.

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*This work is supported in part by funds provided by the U. S. Atomic Energy Commission, the Office of Naval Research, and the Air Force Office of Scientific Research.

†On the staff of the University of California Radiation Laboratory during the performance of this experiment.

[1]Cool, Piccioni, and Clark, Phys. Rev. 103, 1082 (1956).

[2]H. C. Burrowes et al., preceding Letter [Phys. Rev. Lett. 2, 117 (1959)].

[3]DeWire, Jackson, and Littauer, Phys. Rev. 110, 1208 (1958); P. C. Stein and K. C. Rogers, Phys. Rev. 110, 1209 (1958); Heinberg, McClelland, Turkot, Wilson, Woodward, and Zipoy, Phys. Rev. 110, 1209 (1958); Sellen, Cocconi, Cocconi, and Hart, Phys. Rev. 110, 779 (1958).

[4]F. P. Dixon and R. L. Walker, Phys. Rev. Lett. 1, 142 (1958).

[5]R. R. Wilson, Phys. Rev. 110, 1212 (1958).

[6]A. R. Erwin, Jr., and J. K. Kopp, Phys. Rev. 109, 1364 (1958).

π^--p SCATTERING BELOW 1 Bev

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We have carried out measurements of π^--p total and differential elastic cross sections in the energy range 400 to 800 Mev. This work was done by exposing a propane bubble chamber in a pion beam at the Brookhaven Cosmotron.

The chamber had an active volume of 12 in. $\times$6 in. $\times$6 in. which was viewed with two cameras at right angles. The right-angle stereoscopy permits more accurate angle measurements than do more conventional small-angle stereo systems.

Since the intensity of the Cosmotron was not under our control, it was necessary to sacrifice momentum resolution for beam intensity. Therefore, only one analyzing magnet was used which gave pions of momenta 590 ± 50, 750 ± 50, and 890 ± 75 Mev/c. In all, we took about 60 000 pictures.

During the process of scanning, all pion interactions were counted and placed in the following categories: elastic scattering, quasi-elastic scattering, charge exchange, pion production (i.e., apparently in a π^--p collision), diffraction scattering, and carbon stars. All the apparently elastic events were measured and calculated. Because of the right-angle stereo, visual discrimination is usually good and about 70% of the calculated events were consistent with being π^--p elastic scatterings. Thus we have the number of π^--p elastic scatterings and all other events. From this and the total π^--p [1] and π^--C [2,3] cross sections we can deduce the total elastic π^--p cross section. We estimate the π^--p inelastic and π^--p charge exchange from the ratios observed. It is certain that perhaps 30-40% of these inelastic and charge exchange events are the result of collisions with protons bound in carbon. We assume that the ratios on free and bound protons are roughly the same. McCormick and Baggett[4] have measured σ_{el}, σ_{inel}, and σ_{ex} at 810 Mev. They find ratios of these three quantities to be 14.3, 15.5, 6.4, and we find at 760 Mev the ratios 16.7, 13.5, 7.35. Likewise our 460-Mev results have been closely confirmed by unpublished measurements in hydrogen by Stranahan, Ashkin, and DeBenedetti.[5]

The results of our measurements are shown in Figs. 1 and 2. It is perhaps necessary to point out again that the elastic cross sections were measured on the basis of the kinematics alone; that is, on the basis of coplanarity, angle measurement, and, in a considerable fraction of cases, the range of the protons. There is a small background due to quasi-elastic scatterings. We estimate from the angular dispersion of the quasi-elastic events that this background

is less than 5%. Our selection criteria and measurements are essentially the same as those used by Pless[6] who finds a similar background in p-p scatterings in a hydrocarbon chamber at a comparable momentum in a very well resolved beam.

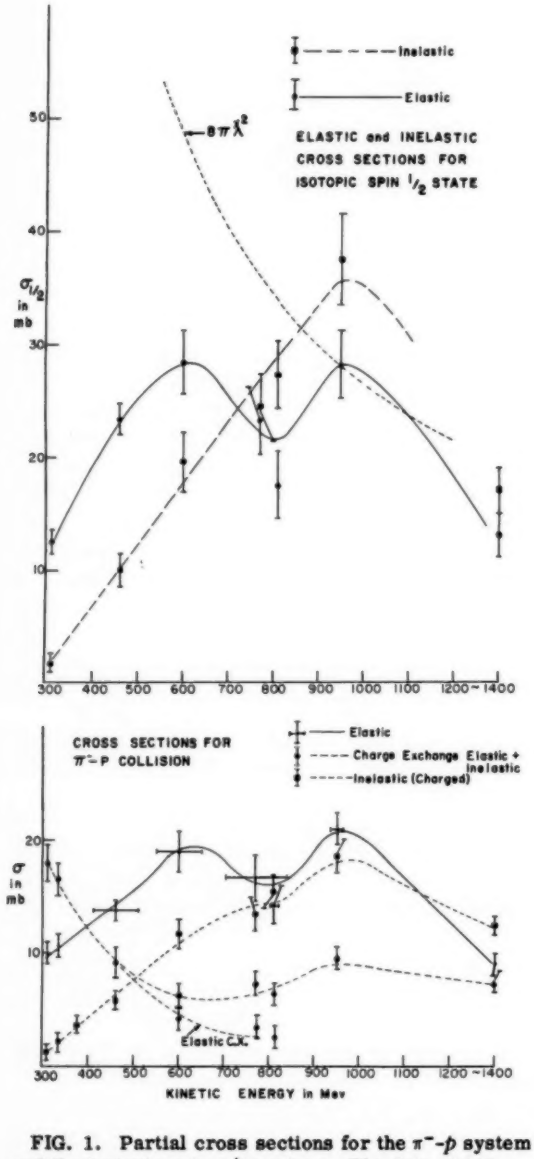

FIG. 1. Partial cross sections for the π^--p system and the isotopic spin 1/2 system. The data at 307, 330, and 375 Mev are the Russian work[7]; at 810 Mev they are from Berkeley[4] and have been normalized to a total cross section of 36 mb. The data at 950 Mev[8,9] and 1.4 Bev are from earlier bubble chamber, diffusion chamber, and emulsion experiments.

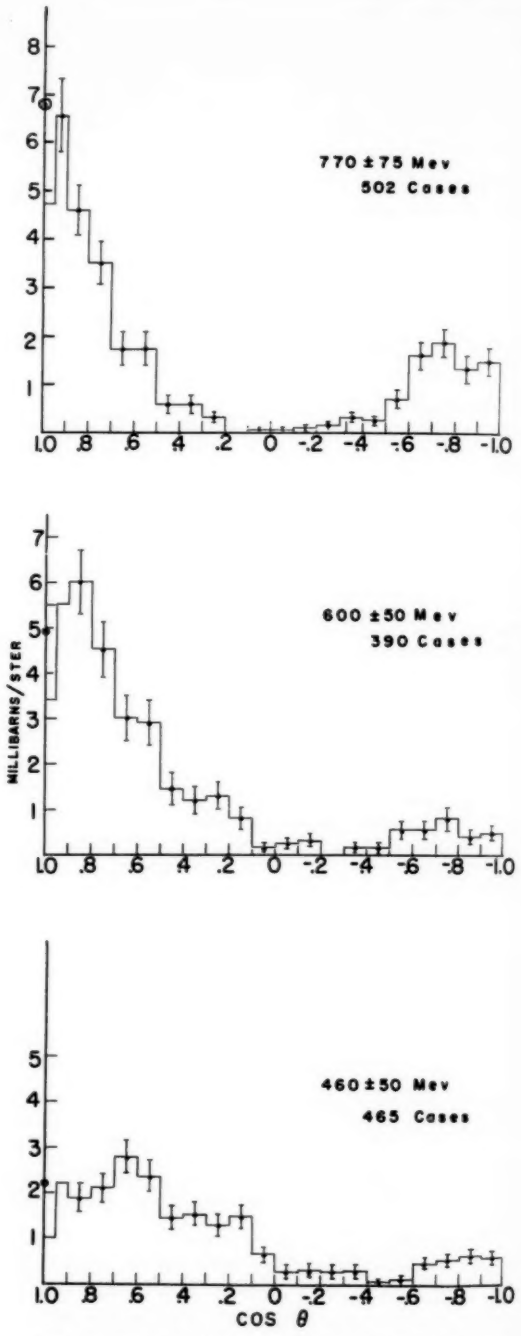

FIG. 2. Differential π^--p elastic scattering cross sections. In most cases the most nearly forward interval (1.0 - 0.95) is not very trustworthy. The circles on the ordinate are the dispersion relation predictions as calculated by Sternheimer from the data of Cool et al.[1]

We give in Fig. 2 our differential angular distributions. It should be noted that at all energies above 300 Mev,[7] the distributions are not apparently dominated by a single J state. At 460 Mev, for example, s, p, and d waves contribute comparable fractions to the elastic cross section. There is quite a marked change in angular distribution between energies of 600 and 760 Mev. A large bump in the backward hemisphere develops. Of the order of 35% of the total elastic cross section is contained in the backward hemisphere. Previously Walker et al.[8] and Erwin and Kopp[9] have commented on the possible origin of the rise as being due to spin-flip scattering. On the basis of the present data we would say that either a $T=1/2$ phase shift has changed sign or goes through 90° in the neighborhood of 600 Mev. It is also possible that the rapid growth of the f-wave phase shift has caused the change.

Recent work on photoproduction of pions in the same energy range at the California Institute of Technology and Cornell University[10] gives indications of a resonant state which should show up around 600 Mev in a π^--p scattering. Our data are certainly consistent with such a picture. The elastic cross section comes up before the inelastic cross section and shows a slight maximum at about 600 Mev. The marked change in shape in the differential elastic cross section between 600 and 760 Mev is likewise consistent with a resonance.

We have extracted the $T=1/2$ cross sections by subtracting out the $T=3/2$ contribution. This is possible because of recent data on π^+-p interactions of Willis et al.[11] and Erwin and Kopp.[12] In order to calculate the $T=1/2$ inelastic cross section it is necessary to know, in addition to the inelastic processes giving charged particles, the inelastic charge exchange cross section, i.e., $\pi^- + p \rightarrow 2\pi^0 + n$. We have measured the amount of $1\pi^0$ and $2\pi^0$ production at 460 and 600 Mev. We did this by counting the number of pairs within four centimeters from the stopping point. This gives the average number of γ rays produced per charge exchange which is simply related to the number of π^0's. At 810 Mev and 760 Mev we use the isobar model,[13] which is consistent with the data, to compute the $2\pi^0$ cross section.

To see in detail which, if any, state resonates at about 600 Mev is a very difficult problem. At 600 Mev, $8\pi\lambda^2$ (the maximum cross section for a $J=3/2$ state) is twice the observed elastic cross section in the $T=1/2$ state. This shows that if the peak in the cross section is produced by a $J=3/2$ wave then the wave must be strongly absorbed. Preliminary efforts at a phase-shift analysis by one of the authors (W.D.W.) supports this view. The phase-shift analysis requires an enhanced $d_{3/2}$ phase shift.[14] A large positive α_{13} seems excluded because of the relatively small amount of charge exchange scattering. The humped appearance of the differential cross section at 460 Mev is characteristic of constructive interference between p- and d-wave spin-flip terms. At 460 Mev preliminary estimates of the $T=1/2$ phase shifts give for the real part of the phase shifts $\alpha_1 = +20°$, $\alpha_{11} = +30°$, $\alpha_{13} = -10°$, $\delta_{13} = +23°$, $\delta_{15} = -5°$. To account for the large peak in the forward direction at 600 Mev, δ_{13} must be considerably increased.

The inelastic cross section in the $T=1/2$ state shows a relatively large peak in the neighborhood of 900 Mev. To account for the size of the cross section requires complete absorption of s, p, and d waves if these are the only waves to be considered. Complete absorption of these waves would be inconsistent with angular distribution of the elastically scattered π's at these energies (in particular the large hump in the backward hemisphere). Consequently it is certain that f waves play an important role in the 800-900 Mev region.

This work has been made possible through the cooperation and help of many people. At Brookhaven we wish to thank Dr. G. B. Collins, Dr. W. Moore, and Dr. R. K. Adair. At Wisconsin the cooperation of the Physics Department shop and in particular P. Schultz and A. Swenson was most helpful. The scanning help of R. Hess is acknowledged. We wish to thank Dr. D. Glaser, Dr. R. Sachs, and Dr. A. Erwin for helpful discussions of various phases of the work. The help of J. Huang in scanning and programming some of the data at Michigan State is gratefully acknowledged. A helpful conversation on photoproduction with K. Berkelman is acknowledged.

*Supported in part by the U. S. Atomic Energy Commission, Wisconsin Alumni Research Foundation.

†Supported in part by a grant from the National Science Foundation.

[1]Cool, Clark, and Piccioni, Phys. Rev. 103, 1082 (1956); also Burroughs, Frisch, Hill, Ritson, Schluter, and Wahlig (to be published). We wish to thank the latter authors for permission to use their results.

[2]Cronin, Cool, and Abashian, Phys. Rev. 107, 1121 (1957).

[3]We wish to thank Dr. Cool and Dr. Cronin for in-

forming us of some of their unpublished cross sections.

[4]L. Baggett, University of California Radiation Laboratory Report UCRL-8302 (unpublished). Also B. McCormick and L. Baggett, in 1958 Annual International Conference on High-Energy Physics at CERN, edited by B. Ferretti (CERN, Geneva, 1958).

[5]S. DeBenedetti (private communication).

[6]I. A. Pless, Phys. Rev. 104, 205 (1956).

[7]S. M. Korenchenko and V. G. Zinov, J. Exptl. Theoret. Phys. (U.S.S.R.) 33, 1307 (1957) [translation: Soviet Phys. JETP 6, 1006 (1958)].

[8]Walker, Hushfar, and Shephard, Phys. Rev. 104, 526 (1956).

[9]A. R. Erwin and J. Kopp, Phys. Rev. 109, 1364 (1958).

[10]Heinberg, McClelland, Turkot, Woodward, Wilson, and Zipoy, Phys. Rev. 110, 1211 (1958); J. I. Vette, Phys. Rev. 111, 622 (1958); F. P. Dixon and R. L. Walker, Phys. Rev. Lett. 1, 4 (1958); see also reports by R. R. Wilson and R. L. Walker at the 1958 Annual International Conference on High-Energy Physics at CERN, edited by B. Ferretti (CERN, Geneva, 1958); Sellen, Cocconi, and Hart, Phys. Rev. 110, 779 (1958).

[11]W. J. Willis, Ph. D. thesis, Yale University, 1958 (unpublished). Also Morris, Rahm, Rau, Thorndike, and Willis, Bull. Am. Phys. Soc. Ser. II, 3, 33 (1958).

[12]A. R. Erwin (private communication).

[13]S. J. Lindenbaum and R. M. Sternheimer, Phys. Rev. 109, 1723 (1958).

[14]The possibility of a $d_{3/2}$ resonance to explain the peak in the photoproduction cross section was expressed in a quantitative way by R. F. Peierls, Phys. Rev. Lett. 1, 174 (1958).

THEORY OF SPIN-ORBIT INTERACTION IN NUCLEAR FORCES*

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It has repeatedly been pointed out in recent years[1-4] that a spin-orbit interaction between two nucleons is necessary to explain the observed scattering of nucleons. From a theoretical point of view the existence of a spin-orbit interaction is not at all surprising, because it has been shown by Breit[5] that in a relativistic treatment of the interaction of nucleons the spin-orbit interaction arises in a natural way. However, pion-theoretical calculations by Klein[6] and several other authors[7] show that the pion theory is unable to account for the large spin-orbit interaction, which is required to explain the experimental results. It is, therefore, necessary to look for some other explanation of the spin-orbit interaction.

Recently we have predicted[8,9] the existence of a hitherto unobserved neutral scalar meson, the ρ^0 meson, which is coupled strongly to the nucleons. Since the mass of the ρ^0 meson is considerably larger than the pion mass, it leads to a force of very short range between the nucleons. The second-order nuclear potential due to the ρ^0 meson is given by

$$V_2(\rho^0) = -\frac{g'^2}{4\pi r} e^{-\lambda' r} + \frac{g'^2}{4\pi r}\frac{1}{2\kappa^2}\frac{d}{dr}\left(\frac{e^{-\lambda' r}}{r}\right)\vec{L}\cdot\vec{S}, \quad (1)$$

where g' is the coupling constant for the interaction of ρ^0 mesons and nucleons, λ' and κ are related to the ρ^0-meson mass μ' and the nucleon mass M as $\lambda' = \mu' c/\hbar$ and $\kappa = Mc/\hbar$, and we have used the Signell-Marshak definitions[3] of $\vec{L}$ and $\vec{S}$. The coefficient of $\vec{L}\cdot\vec{S}$ in (1) can be expressed as

$$V_{LS} = \frac{V^0}{x}\frac{d}{dx}\left(\frac{e^{-nx}}{x}\right), \quad (2)$$

with

$$V_0 = \frac{g'^2}{4\pi c\hbar}\frac{\lambda c\hbar}{2}\left(\frac{\lambda}{\kappa}\right)^2, \quad (3)$$

where λ is related to the pion mass μ as $\lambda = \mu c/\hbar$, while $x = \lambda r$ and $n = \mu'/\mu$.

According to our earlier ideas,[8] the ρ^0-meson mass should be somewhat larger than twice the pion mass, and the coupling constant for the interaction of ρ^0 mesons and nucleons should have the same value as the coupling constant for pions and nucleons. Thus, we can take

$$n \approx 2, \quad g'^2/4\pi c\hbar \approx 14. \quad (4)$$

We also have

$$\lambda/\kappa = 1/6.7, \quad \lambda c\hbar = \mu c^2 = 139.4 \text{ Mev}, \quad (5)$$

where we have taken the pion mass as $273 m_e$. Substituting the above values in (2) and (3), we find

$$V_{LS} = \frac{V_0}{x}\frac{d}{dx}\left(\frac{e^{-2x}}{x}\right), \quad (6)$$

with

$$V_0 = 21.7 \text{ Mev}. \quad (7)$$

It seems to us quite astonishing that not only (6) has exactly the same form as the latest phenomenological spin-orbit interaction of Signell, Zinn, and Marshak,[4] but our theoretical value of V_0 is also remarkably close to the phenomeno-

logical value of $V_0 = 21$ Mev.

It is tempting to conclude that the ρ^0 meson provides a complete explanation of the spin-orbit interaction between two nucleons. The resulting interaction also seems to be in agreement with the requirements of the shell model of the heavier nuclei.[10] It must, however, be noted that some objections have also been raised[11] against isotopic-spin-independent spin-orbit interactions, and at the present stage of our knowledge we cannot settle the problem of spin-orbit interaction with complete certainty. It is also possible that the ρ^0 meson is coupled less strongly to the nucleons than the pions, and that the spin-orbit interaction is partly due to the ρ^0 meson and partly due to pions.

*Supported in part by the National Science Foundation.

[1]K. M. Case and A. Pais, Phys. Rev. 80, 203 (1950).
[2]J. L. Gammel and R. M. Thaler, Phys. Rev. 107, 291, 1337 (1957).
[3]P. S. Signell and R. Marshak, Phys. Rev. 106, 832 (1957); 109, 1229 (1958).
[4]Signell, Zinn, and Marshak, Phys. Rev. Lett. 1, 416 (1958).
[5]G. Breit, Phys. Rev. 51, 248, 778 (1937); 53, 153 (1938).
[6]A. Klein, Phys. Rev. 90, 1101 (1953).
[7]For the latest work on this subject, see G. Breit, Phys. Rev. 111, 652 (1958) and S. Otsuki, Progr. Theoret. Phys. (Japan) 20, 171 (1958).
[8]S. N. Gupta, Phys. Rev. 111, 1436 (1958).
[9]S. N. Gupta, Phys. Rev. 111, 1698 (1958).
[10]J. P. Elliott and A. M. Lane, Phys. Rev. 96, 1160 (1954).
[11]H. Feshbach, Phys. Rev. 107, 1626 (1957); A. M. Sessler and H. M. Foley, Phys. Rev. 110, 995 (1958).

V-A FOUR-FERMION INTERACTION AND THE INTERMEDIATE CHARGED VECTOR MESON*

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The universal *V-A* form of Fermi interaction seems to be compatible with the current experimental results.[1,2] In order to forbid many unwanted reactions, it has been speculated that the weak Fermi interactions are realized through the interaction of a charged current, J_α, with itself by exchange of a heavy charged boson.[1] The simplest possibility is the introduction of a charged vector meson B_α (with mass m_B) which is coupled to the J_α with the coupling constant F:

$$F J_\alpha B_\alpha + \text{H.c.} \qquad (1)$$

The J_α may consist of the lepton currents, $\bar{e}\gamma_\alpha \times(1+\gamma_5)\nu$ and $\bar{\mu}\gamma_\alpha(1+\gamma_5)\omega$, and of the strangeness conserving and nonconserving baryon currents. We deliberately denote the neutral counterpart of the muon as ω. First of all, it may be remarked that the B_α-meson must be heavier than the K-meson.[3] The case when the B_α-meson is extremely heavy may not be realistic and there may be little point in introducing it. One interesting theoretical evidence against a B_α-meson which is not unreasonably heavy has been pointed out. Namely, if we take the two-component theory of the neutrino ($\nu \equiv \omega$) with lepton number conservation, the existence of the B_α would cause the $\mu \to e + \gamma$ transition.[4] It is, however, to be remembered that we could forbid this decay if we do not assume $\nu \equiv \omega$. For instance,[5] one may interchange the lepton number of μ^+ and μ^- and at the same time change the neutrino accompanying μ to an antineutrino[6] ($\omega \equiv \nu^C$). This possibility can hardly be differentiated from the usual theory by presently feasible experiments. In view of this situation, it may be worthwhile to investigate further the possible existence of a B_α-meson. In the following, unless mentioned, we neglect electromagnetic corrections. For the $\pi \to \mu(e) + \nu$ and $K \to \mu(e) + \nu$ decays nothing would be changed except for the replacement of the usual Fermi coupling constant G by $\sqrt{2}\, F^2/m_B{}^2$. For the K_{e3} and $K_{\mu 3}$ decays, however, the effect may be observable by future experiments. For instance, the general form of the matrix element of the K_{e3} mode is given (neglecting the electron mass) by

$$M F^2 \bar{e}(k\cdot\gamma)(1+\gamma_5)\nu(A - 2m_K E_\pi)^{-1}\phi_K(k)\,\phi_\pi(k-p), \qquad (2)$$

where $A = m_B{}^2 - m_K{}^2 - m_\pi{}^2$. k_α and $k_\alpha - p_\alpha$ denote the energy-momentum four-vector of the K-meson and the pion, respectively. In principle, the effective coupling parameter M may depend on the total pion energy E_π. As M represents the contribution of baryon-antibaryon loops, this dependence is probably rather weak.[7] The denominator of (2) expresses the propagator of the B_α-meson. If the mass of the B_α-meson is not far removed from the K-meson mass, this denominator will behave as $-2m_K E_\pi$ [$A = 0$ if $m_B = (m_K{}^2 + m_\pi{}^2)^{1/2}$]. In this case the effect may be observ-

able particularly for the pion spectrum. In fact, the pion energy spectrum $W(E_\pi)dE_\pi$ given by the effective interaction (2) will have the proportionality

$$W(E_\pi) \propto \frac{(E_\pi^2 - m_\pi^2)^{3/2}}{(A - 2m_K E_\pi)^2} \approx \frac{1}{4m_K^2} \frac{(E_\pi^2 - m_\pi^2)^{3/2}}{E_\pi^2}, \quad (3)$$

whereas the contact V-A interaction gives[8]

$$W(E_\pi) \propto (E_\pi^2 - m_\pi^2)^{3/2}. \quad (4)$$

These spectra are shown in Fig. 1. The electron spectrum given by (2) is shown in Fig. 2. For kinematical reasons the electron spectrum does not show a marked difference from the ordinary case.[9] At present the experimental material is still too scanty to permit any conclusion to be drawn.[10] However, we may hope that an analysis

FIG. 1. Pion energy spectra of the K_{e3} decay given by (3) and (4). They are normalized to give the same area.

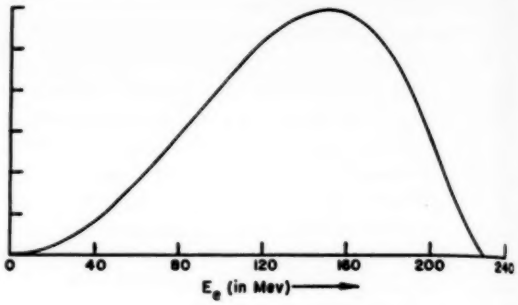

FIG. 2. Electron energy spectrum of the K_{e3} decay in the present model. Compare with references 8 and 9.

of a large number of K_{e3}^0 as well as K_{e3}^+ decays should give more information. We have neglected the possible dependence of the parameter M on E_π. With this approximation the energy spectra of $K^\pm \to e^\pm + \nu + \pi^0$, $K_2^0 \to e^\pm + \nu + \pi^\mp$, and $K_1^0 \to e^\pm + \nu + \pi^\mp$ would be the same except for the small effect of the mass difference of the participating particles. In general, if the dependence of M on E_π is marked, these spectra may be different from each other.[11] It may be noted that the experiment in which E_π is fixed cannot be used for our purpose.

Finally, we would add some comments on the μ-decay. The effective matrix element would take the form (putting $m_e = 0$)

$$(F^2/m_B^2)\,\bar{\omega}\gamma_\alpha(1+\gamma_5)\mu(1+k^2/m_B^2)^{-1}\bar{e}\gamma_\alpha(1+\gamma_5)\nu, \quad (5)$$

where k_α is the energy-momentum transfer between two pairs (μ, ω) and (e, ν). The Michel parameter of the interaction (5) is approximately given[12] by $\rho = 0.75 + \frac{1}{3}(m_\mu/m_B)^2$. The recently reported[13] value of ρ (corrected for the electromagnetic corrections) is 0.79 ± 0.03. If we take $m_B \approx m_K$, the nonlocal effect would be ≈ 0.015 which is not in contradiction with the above experiment. Furthermore, the decay rate computed from (5) will be $W = W_0\{1 + \frac{3}{5}(m_\mu/m_B)^2\}$ where W_0 is the rate for the contact interaction. The universal coupling constant G of the contact Fermi interaction estimated from O^{14} decay (after inclusion of the electromagnetic corrections) is given by[14] $G = (1.37 \pm 0.02) \times 10^{-49}$ erg cm^3. Then the predicted value of the muon mean life corrected for the electromagnetic effects turns out to be $(2.33 \pm 0.05) \times 10^{-6}$ sec, while the experimental value is $(2.22 \pm 0.02) \times 10^{-6}$ sec. Now for the β-decay the nonlocal effect is completely negligible. If we neglect the change of

the electromagnetic corrections due to the finiteness of m_B, then the above-discussed nonlocal effect decreases the value $(2.33 \pm 0.05) \times 10^{-8}$ sec to $(2.27 \pm 0.049) \times 10^{-8}$ sec which is in the right direction to maintain the universality of the coupling constant. It is, however, necessary to treat the radiative corrections correctly (without assuming $m_B = \infty$) in order to get more precise information. This may be discussed elsewhere. We would like to thank Dr. J. Sucher for his interest in this work.

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†On leave of absence from the Institute of Theoretical Physics, Kanazawa University, Japan.

[1]R. P. Feynman and M. Gell-Mann, Phys. Rev. 109, 193 (1958).

[2]E. C. G. Sudarshan and R. E. Marshak, Proceedings of the Padua-Venice Conference, September, 1957 [Suppl. Nuovo cimento (to be published)], and Phys. Rev. 109, 1860 (1958).

[3]The existence of the B_α-meson will give severe restrictions on the mass spectra of observable particles.

[4]For instance, G. Feinberg, Phys. Rev. 110, 1482 (1958).

[5]J. Schwinger, Ann. Phys. 2, 407 (1957); K. Nishijima, Phys. Rev. 108, 907 (1958); M. Konuma, Nuclear Phys. 5, 504 (1958). It may be remarked that this procedure alone is not enough to forbid $K^+ \to \pi^+ + \nu + \bar{\nu}$ and $K^+ \to \pi^+ + e^+ + e^-$, etc.

[6]In the case $\omega \equiv \nu^{\mathcal{C}}$ the neutrinos have four components rather than two. This is a modified form of the twin neutrino theory where one twin is associated with the e^- and the other twin with the μ^+. The complete Lagrangian is invariant under $\nu \to \gamma_5 \nu$, $e \to e$ and $\mu \to -\mu$ (Nishijima transformation[5]) which guarantees the strict vanishing of the neutrino rest mass.

[7]We could expand M into $a + b\, k \cdot p/m^2 + \ldots$, where m is a quantity of the order of baryon mass.

[8]S. Oneda, Nuclear Phys. 4, 21 (1957), Eq. (5); S. W. Macdowell, Nuovo cimento 6, 1445 (1957); S. Furuichi, Nuovo cimento 7, 269 (1958).

[9]Furuichi, Kodama, Ogawa, Sugawara, Wakasa, and Yonezawa, Progr. Theoret. Phys. (Japan) 17, 89 (1957).

[10]Bruin, Holthuizen, and Jongejans, Nuovo cimento 9, 422 (1958). Earlier references will be found there.

[11]However, if, for instance, the strangeness-nonconserving baryon current behaves as an isospinor ($|\Delta I| = 1/2$), these energy spectra should be the same irrespective of the dependence of M on E_π. Okubo, Marshak, Sudarshan, Teutsch, and Weinberg, Phys. Rev. 112, 665 (1958). If the current is a mixture of $|\Delta I| = 1/2$ and $3/2$, this will in general not be the case.

[12]T. D. Lee and C. N. Yang, Phys. Rev. 108, 1611 (1957).

[13]R. J. Plano and A. Lecourtois, report at the Conference on Weak Interactions, Gatlinburg, Tennessee, October, 1958 [Bull. Am. Phys. Soc. 4, 82 (1959)].

[14]For instance, S. M. Berman, Phys. Rev. 112, 267 (1958); T. Kinoshita and A. Sirlin (to be published).

γ_5-SYMMETRIZATION

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By imposing the requirement that a Fermi interaction Lagrangian $(\bar{\psi}_1 \Omega \psi_2)(\bar{\psi}_3 \Omega \psi_4)$ remain invariant for the independent transformations $\psi_i \to \gamma_5 \psi_i$ ($m_i \to -m_i$,[1] $i = 1, 2, 3, 4$), Sakurai[2] was able to deduce that weak interactions must be of (V-A) type. Unfortunately (the parity-conserving) strong interactions are not invariant for this transformation. In this connection the following remarks may be made:

(1) The Yukawa interactions

$$\bar{\psi}\psi\phi, \quad \bar{\psi}\gamma_5\psi\phi, \quad \bar{\psi}\gamma_\mu\psi A_\mu, \quad \bar{\psi}\gamma_\mu\gamma_5\psi A_\mu \qquad (1)$$

remain unchanged for the following substitutions: $\psi \to \gamma_5 \psi$, $\phi \to -\phi$, $A_\mu \to A_\mu$, $m \to -m$. Alternatively stated, all Lagrangians of the above type involving one ψ field are γ_5-symmetrized in the following sense:

$$L = \tfrac{1}{2}[\mathfrak{L}(\psi, A_\mu, \phi, m) + \mathfrak{L}(\gamma_5\psi, A_\mu, -\phi, -m)]. \qquad (2)$$

(2) If there are two independent Fermi fields ψ_1 and ψ_2 and all interactions are of Yukawa type, define a γ_5-symmetrized Lagrangian as

$$\begin{aligned} L = \tfrac{1}{4}[&\mathfrak{L}(\psi_1, \psi_2, A_\mu, \phi, m) + \mathfrak{L}(\gamma_5\psi_1, \psi_2, A_\mu, -\phi, -m) \\ &+ \mathfrak{L}(\psi_1, \gamma_5\psi_2, A_\mu, -\phi, -m) \\ &+ \mathfrak{L}(\gamma_5\psi_1, \gamma_5\psi_2, A_\mu, -\phi, -m)]; \end{aligned} \qquad (3)$$

i.e., change every field ψ independently into $\gamma_5\psi$ and after making the appropriate changes of ϕ fields (and m) sum all possible terms. The generalization to the case of n different Fermi fields is immediate and involves symmetrization of 2^n terms.

It may easily be checked that the only γ_5-symmetrized interaction terms which conserve parity are

(a) nonderivative interactions of scalar bosons, and

(b) interactions of vector mesons when only one

Fermi field is involved (e.g. $\bar{\psi}\gamma_\mu\psi A_\mu$). γ_5-symmetrized vector meson interactions involving two different Fermi fields must exist in $V \pm A$ combinations. It may be noted that this symmetrization commutes with any isotopic symmetries Lagrangians may possess.

It has been suggested that weak[3] interactions are mediated by charged vector bosons. It is to be remarked that, if this hypothesis is accepted, the complete strong, electromagnetic, and weak interaction, which is believed to describe the actual physical world, is γ_5-symmetrized in the sense used above. Alternatively the prescription of γ_5-symmetrization with Yukawa interactions could have been used to predict almost uniquely the presently accepted physical Lagrangian.

[1] $m_i \rightarrow -m_i$ is necessary for the invariance of the corresponding free Lagrangian.

[2] J. J. Sakurai, Nuovo cimento 7, 649 (1958). The formalism of R. P. Feynman and M. Gell-Mann [Phys. Rev. 109, 193 (1958)] and E. C. G. Sudarshan and R. E. Marshak [Phys. Rev. 109, 1860 (1958)] gives identical results. However there is one difference: in Sukarai's work there is no insistence on the projection operators $(1+\gamma_5)$ and $(1-\gamma_5)$ to start with. These particular combinations of operators arise naturally where required during the course of theoretical development.

[3] R. P. Feynman and M. Gell-Mann (reference 2).

ABSTRACTS

In this section are printed the abstracts of Articles that have been forwarded to The American Institute of Physics for publication in THE PHYSICAL REVIEW. In quoting information obtained from this section before the appearance of the corresponding Article, reference should be made to "Physical Review (to be published)" rather than to this Journal.

THEORY OF MANY-BOSON SYSTEMS. PAIR THEORY. M. Girardeau* and R. Arnowitt, Department of Physics, Syracuse University, Syracuse, New York (Received September 11, 1958).

The many-boson system with repulsive interactions is treated by a variational method based on a variational trial state of an exponential pair-excitation type obtained by a generalization of that of Bogoliubov; the treatment is closely related to an intermediate-coupling approximation with respect to pairs. The nonlinear integral equation which determines the variational ground state is derived, and various properties of this ground state and the associated excited states are examined. The resultant low-lying spectrum lies below Bogoliubov's by an amount proportional to the total number of particles. The variational principle is shown to produce rigorous energy eigenvalue differences for the pair part of the Hamiltonian. The variational states, however, still exhibit unphysical features characteristic of pair-excitation states: the pair correlation function does not go to zero at zero particle separation and the phonon spectrum exhibits a gap above the ground state. It is suggested that these features can be removed by using states which take into account excitation of momentum-conserving groups of more than two particles.

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BINARY COLLISION METHOD IN QUANTUM STATISTICAL MECHANICS. I. GENERAL FORMULATION. T. D. Lee, Columbia University, New York, New York, and C. N. Yang, Institute for Advanced Study, Princeton, New Jersey (Received October 2, 1958).

The general formulation of the binary collision method is presented. It is shown that the grand partition function of an N-body system can be evaluated by using a binary kernel which is itself obtainable from a solution of the two-body problem. The limit for an infinite system and the case of a singular repulsive potential are discussed.

THEORY OF THE THERMAL CONDUCTIVITY OF SUPERCONDUCTORS. J. Bardeen, G. Rickayzen, and L. Tewordt, Physics Department, University of Illinois, Urbana, Illinois (Received October 8, 1958).

A theory of the thermal conductivity of superconductors is presented, based on the theory of superconductivity due to Bardeen, Cooper, and Schrieffer. The excited states of the system are treated as quasi-particles allowing a Boltzmann equation to be set up. The electronic contribution to the thermal conductivity when the dominant scatterers are impurities has been calculated exactly. The result is very close to that of the Heisenberg-Koppe theory which is in fair agreement with experiment. The variational principle of Wilson has been used to find the electronic conductivity when the dominant scatterers are lattice waves. It is concluded that the theory fails to predict the sharp drop in the ratio κ_{es}/κ_{en} as the temperature is lowered below T_c, a feature which is characteristic of the experimental results. The effect of the electrons on the lattice conductivity has also been calculated. The theoretical values may be too large.

ENERGY LEVELS IN IRRADIATED GERMANIUM. E. I. Blount, Westinghouse Research Laboratories, Pittsburgh, Pennsylvania (Received October 2, 1958).

The energy levels found in germanium irradiated by different particles seem at first to be mutually inconsistent. It is possible tentatively to reconcile the differences by consideration of clustering and association of defects. Four levels are ascribed to single vacancies and interstitials. A crude theory is constructed to explain these levels, particularly their asymmetrical distribution in the energy gap, and to assign each to a definite defect. This theory differs somewhat from a previous one due to James and Lark-Horovitz; some differences in experimental predictions are discussed in particular.

TWO ELECTRICAL PHENOMENA ON LIQUID GERMANIUM SURFACES. A. I. Bennett, Westinghouse Research Laboratories, Pittsburgh, Pennsylvania (Received October 8, 1958).

Part I of this paper describes some experiments in which an electric discharge, between the surface of a molten pool of germanium and a nearby solid germanium electrode, was found to repel surface scum from the vicinity of the discharge. Part II reports some experiments in which it was found possible to produce crystal nucleation by the application of a relatively low electric field to the surface of a supercooled germanium melt. No explanation for either of these phenomena is proposed.

HEAT CAPACITY OF PSEUDO-COPPER BELOW 4.2°K. G. L. Guthrie, John Jay Hopkins Laboratory for Pure and Applied Science, General Atomic Division of General Dynamics Corporation, San Diego, California (Received September 23, 1958).

Heat-capacity measurements have been made on two face-centered cubic alloys of the type $Cu_{1-2x}Zn_xNi_x$ at temperatures below 4.2°K for the values $x \approx 0.03$ and $x \approx 0.08$. The heat capacity has been found to obey a law of the type $C_v = \gamma T + (464.5/\theta_D^3)T^3$ cal/mole (°K). The Debye theta is nearly the same as that for pure copper. For an alloy having $x \approx 0.03$, θ_D is approximately 1% higher than θ_D for pure copper. However, γ increases rapidly with increasing x, and is approximately 20% higher than the γ for pure copper when $x \approx 0.03$. The results are found to be incompatible with a rigid-band model approach to alloy theory.

MODEL FOR LATTICE THERMAL CONDUCTIVITY AT LOW TEMPERATURES. Joseph Callaway, Westinghouse Research Laboratories, Pittsburgh, Pennsylvania (Received October 10, 1958).

A phenomenological model is developed to facilitate calculation of lattice thermal conductivities at low temperatures. It is assumed that the phonon scattering processes can be represented by frequency-dependent relaxation times. Isotropy and absence of dispersion in the crystal vibration spectrum are assumed. No distinction is made between longitudinal and transverse phonons. The assumed scattering mechanisms are (1) point impurities (isotopes), (2) normal three-phonon processes, (3) umklapp processes, and (4) boundary scattering. A special investigation is made of the role of the normal processes which conserve the total crystal momentum and a formula is derived from the Boltzmann equation which gives their contribution to the conductivity. The relaxation time for the normal three-phonon processes is taken to be that calculated by Herring for longitudinal modes in cubic materials. The model predicts for germanium a thermal conductivity roughly proportional to $T^{-3/2}$ in normal material, but proportional to T^{-2} in single-isotope material in the temperature range 50° - 100°K. Magnitudes of the relaxation times are estimated from the experimental data. The thermal conductivity of germanium is calculated by numerical integration for the temperature range 2° - 100°K. The results are in reasonably good agreement with the experimental results for normal and for single-isotope material.

PHOSPHOR WITH FLUORESCENCE LARGER THAN THE ENERGY GAP. Seymour P. Keller and George D. Pettit, Research Laboratory, International Business Machines Corporation, Poughkeepsie, New York (Received October 10, 1958).

Under certain conditions of excitation, a Pr^{+3}-activated SrS phosphor exhibits an emission which is energetically greater than the energy gap. This emission is observed at 77°K and not at room temperature. It is proposed that the excitation and short-wavelength emission take place between levels of the Pr^{+3} center. At 77°K the reduction of lattice vibrations effects the isolation of the Pr^{+3} site from the lattice and hence the probability of hole transfer from the activator to the lattice is diminished. The nature of the energy levels of the Pr ion is discussed.

DEPENDENCE OF GEOMETRIC MAGNETIC ANISOTROPY IN THIN IRON FILMS. T. G. Knorr* and R. W. Hoffman, Case Institute of Technology, Cleveland, Ohio (Received October 7, 1958).

Thin iron films produced by vacuum deposition are observed to display a magnetic anisotropy dependent on the geometric location of the evap-

orating filament. A fiber axis structure is induced during the evaporation. If the direction of incidence of the metallic flux varies from a normal to the substrate this fiber axis tilts in a similar direction. As soon as the fiber axis is no longer normal to the plane of the film, magnetic anisotropy is produced. The variation of the tilt of the fiber axis correlates with and can be shown to account for the geometric dependence of the observed anisotropy. Calculations of the magnetic anisotropy expected in a film with a fiber axis structure and an anisotropic tensile stress in the plane of the film agree with the experimental observations for those films which show only rotational switching. Anisotropies of 3×10^4 ergs/cm^3 have been observed in iron films about 350 A thick prepared without the use of applied magnetic fields either during deposition or during a subsequent thermal anneal.

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TEMPERATURE DEPENDENCE OF ANISOTROPY ENERGY IN ANTIFERROMAGNETS. P. Pincus, Department of Physics, University of California, Berkeley, California (Received October 7, 1958).

The Akulov-Zener classical theory of the temperature dependence of ferromagnetic anisotropy energy is extended to the antiferromagnetic case. The result, for short-range interactions only, is that

$$K_n(T)/K_n(0) = [M(T)/M(0)]^{n(n+1)/2},$$

where $K_n(T)$ is the nth order anisotropy constant and $M(T)$ is the sublattice magnetization. Here, $K_n(0)$ and $M(0)$ refer to the corresponding quantities at absolute zero. The result is verified by a spin wave calculation.

CALCULATION OF MIGRATION AND BINDING ENERGIES OF MONO-, DI-, AND TRIVACANCIES IN COPPER WITH THE USE OF A MORSE FUNCTION. A. C. Damask* and G. J. Dienes, Department of Physics, Brookhaven National Laboratory, Upton, New York, and V. G. Weizer, Lewis Research Center of the National Aeronautics and Space Administration, Cleveland, Ohio (Received October 7, 1958).

A machine calculation has been made of the migration and binding energy of a trivacancy in copper with the use of a Morse function. It is found that a very large relaxation occurs for one atom into the trivacancy. This relaxation causes the trivacancy to be shared equally by four atomic sites and results in a large contribution (about 2.3 ev) to the binding energy. The migration of a trivacancy requires a partial dissociation of this configuration. The energy of migration is calculated to be 1.9 ev. Thus, a trivacancy is highly stable, quite immobile, and is, therefore, probably the nucleus for void formation. For purposes of comparison the migration energies of mono- and divacancies were also computed by the same method without relaxation. These energies are 1.3 and 0.2 ev, respectively.

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INFLUENCE OF SOLUTES ON SELF-DIFFUSION IN THE FACE-CENTERED CUBIC LATTICE. Howard Reiss, Bell Telephone Laboratories, Murray Hill, New Jersey (Received July 16, 1958; revised manuscript received November 13, 1958).

A theory is given for the influence of substitutional solutes on self-diffusion in the face-centered cubic lattice. The theory is limited to cases in which the concentration of solute is low enough so that only one solute atom at a time can interact with a given tracer atom.

Two different kinds of approximation are employed, one in which the processes of association and dissociation of vacancies and solute atoms do not themselves contribute to transport, and one in which they do but the frequency of exchange between solute and vacancy is considered to be infinite.

From data on the diffusion coefficient of the solute as well as on the self-diffusion coefficient in its dependence on solute concentration, the ratio of the frequency with which a vacancy exchanges with a solute atom to that with which it exchanges with a host atom in the first coordination shell of a solute can be estimated. This ratio appears to lie between 0.1 and 0.5 for solutes in silver which increase self-diffusion and for which experimental data are available.

An analysis is given which shows that a good estimate of the influence of a given solute on self-diffusion can be made when only the diffusion coefficient of that solute is known.

Finally the effect which Pd in silver has on the

self-diffusion coefficient (Pd reduces the self-diffusion coefficient) is calculated on the basis of the theory. Agreement between theory and experiment is satisfactory.

TRANSFORMATIONS OF ISING MODELS. Michael E. Fisher, Wheatstone Physics Laboratory, King's College, London, England (Received August 29, 1958).

The "star-triangle" and "decoration" transformations are generalized so as to apply to arbitrary mechanical systems coupled to the spins of a standard Ising net. This leads to exact solutions for further plane Ising nets and also for lattices in which the spins on alternate sites have a magnitude greater than $S=\frac{1}{2}$. A general class of antiferromagnetic Ising models is constructed; exact closed expressions can be derived for all the thermodynamic and magnetic properties of these models is an arbitrary magnetic field.

The magnetizations and susceptibilities of Ising nets in which different spins have different magnetic moments are investigated and a valuable relation between the susceptibilities of the honeycomb and triangular lattices is derived. It is shown how correlation functions involving a given spin can be expressed in terms of correlations involving the nearest-neighbor spins instead.

ENERGY DISTRIBUTION OF NEUTRONS SCATTERED BY LIQUID LEAD. I. Pelah, W. L. Whittemore,* and A. W. McReynolds, John Jay Hopkins Laboratory for Pure and Applied Science, General Atomic Division of General Dynamics Corporation, San Diego, California (Received October 10, 1958).

An experiment has been performed to test some of the predictions of a recent theory of neutron scattering by liquids proposed by Vineyard. Neutrons of 0.0685 ev were scattered at 90° by lead at 250°, 330°, and 550°C, and the widths at half maximum of the energy distributions were determined. The width at 250°C for solid lead was obtained for comparison purposes and found to be narrower than would be predicted for a perfect-gas model. The results for liquid lead show (1) that the simple diffusion model gives a slightly broader distribution than is observed for neutrons scattered by lead near the melting point and (2) that this model predicts a much larger variation of the width with temperature than is observed experimentally.

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EFFECT OF THERMAL PRETREATMENT ON THE THERMOLUMINESCENCE OF KCl CRYSTALS. A. Halperin and M. Schlesinger, Department of Physics, The Hebrew University, Jerusalem, Israel (Received July 29, 1958).

Thermal pretreatment was found to enhance the thermoluminescence of x-colored KCl crystals. The effect was found to penetrate into the crystal by diffusion from the outer surfaces. It is affected by the atmosphere surrounding the heated crystal, and it seems also to be structure sensitive. For some crystals the enhancement in the thermoluminescence was by a factor as high as 4000 when the number of quanta emitted during the thermoluminescence was about the same as the number of F-electrons bleached during this process.

ELECTRON SPIN RESONANCE IN NICKEL-DOPED GERMANIUM. G. W. Ludwig and H. H. Woodbury, General Electric Research Laboratory, Schenectady, New York (Received October 10, 1958).

Electron spin resonance absorption proportional in intensity to the Ni^- concentration has been detected at 14 kMc/sec in germanium crystals. The spectrum, which consists of six main lines having anisotropic g-values, is interpreted in terms of a Jahn-Teller distortion of the nickel ion in any cubic direction. The principal axes of the g-tensor are the cubic direction and two mutually perpendicular [110] directions having g-values of 2.0294, 2.0176, and 2.1128, respectively. Hyperfine structure due to Ni^{61} in enriched samples shows the same principal axes and principal values of ≤ 1.6, 12.2 and 10.3, respectively, in units of 10^{-4} cm^{-1}, and establishes the spin of that isotope as 3/2. Hyperfine interaction with Ge^{73} in the two nearest neighbor positions toward which the nickel atom moves in the Jahn-Teller distortion has been resolved and ranges from 19.8×10^{-4} cm^{-1} for the magnetic field in the Ni-Ge^{73} direction to 16.5×10^{-4} cm^{-1} at right angles to this axis. Above 20.4°K the Jahn-Teller distortion reorients with increasing rapidity.

The line broadening and eventual merging has been studied and estimates have been made of the activation energies for reorientation by inversion and by rotation.

SURFACE PROPERTIES OF ETCHED TUNGSTEN SINGLE CRYSTALS. F. L. Hughes, H. Levinstein, and R. Kaplan,* Physics Department, Syracuse University, Syracuse, New York (Received June 13, 1958).

An examination of the surface properties of chemically etched W single crystals has been made to determine the possibility of developing a macroscopic surface with adsorption properties equivalent to an ideal crystal plane.

Crystals up to three inches in length have been grown in 0.003-inch W wires by a thermal gradient technique. The crystals completely filled the wire cross section.

The etching appeared to have produced a surface which was homogeneous on a microscopic level. Thermionic emission was observed to be identical for the three crystallographic directions, [112], [11$\bar{2}$], and [110], with a value of $\phi_0^* = 5.25 \pm 0.05$ v. Surface ionization measurements gave the same value of work function within experimental error.

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MEAN ADSORPTION LIFETIME OF Rb ON ETCHED TUNGSTEN SINGLE CRYSTALS. I. IONS. F. L. Hughes and H. Levinstein, Physics Department, Syracuse University, Syracuse, New York (Received July 14, 1958).

Using an interrupted atomic beam technique, the mean adsorption lifetime, τ_i, has been measured for Rb ions adsorbed on an "essentially clean," etched tungsten surface. Etching produced a homogeneous surface with a uniform work function. An estimated 90% of the surface consisted of microcrystal facets of the same crystallographic plane. The maximum Rb concentration was about 10^{10} atoms/cm^2.

The data may be represented by the equation $\tau_i = \tau_i^0 \exp(Q_i/kT)$, in the region $1100 < T < 1300$°K. The weakly temperature-dependent term τ_i^0 appears to be affected by the shape, near the minima, of the ionic potential barrier of height Q_i.

Lifetime measurements indicated that τ_i^0 and Q_i are strongly affected by the physical state of the metal and by the type and degree of contamination. It was found that τ_i^0 as well as Q_i must be used in characterizing "chemical" adsorption processes.

MEAN ADSORPTION LIFETIME OF Rb ON ETCHED TUNGSTEN SINGLE CRYSTALS. II. NEUTRALS. F. L. Hughes, Physics Department, Syracuse University, Syracuse, New York (Received July 14, 1958).

The mean adsorption lifetime has been measured for the fraction of adsorbed Rb which was desorbed as neutral atoms from an etched W surface. The lifetime was found to be greatly different from that obtained in the case of Rb ions. Experimental values have been found to be consistent with a theoretical model which assumes that, in a surface ionization process, the adsorbed particles may occupy either of two discrete states with distinctly different types of binding.

ELECTRON IMPACT IONIZATION OF Ne, O, AND N. M. J. Seaton, Department of Physics, University College, London, England (Received October 13, 1958).

The Bethe approximation gives a functional relation between the cross sections for electron-impact ionization (Q) and photo-ionization (a). Estimates of Q_N and Q_O are obtained using experimental values for Q_{Ne} and calculated values for a_N, a_O, and a_{Ne}.

IONIZATION OF ATOMIC OXYGEN ON ELECTRON IMPACT. Wade L. Fite and R. T. Brackmann, John Jay Hopkins Laboratory for Pure and Applied Science, General Atomic Division of General Dynamics Corporation, San Diego, California (Received October 13, 1958).

The cross section for ionization of atomic oxygen has been measured using modulated atomic beam techniques. First the ratio of the cross sections for production of the molecular oxygen ion and for total ion production in collisions of electrons with oxygen molecules was measured. Then the ratio of the ionization cross section of the free oxygen atom and the cross section for production of the molecular ion in electron-molecule collisions was determined. From the

previously known total ionization cross section of the molecule, and the measured ratios, the unknown cross sections were determined. The experimental results are compared with the calculations of Seaton.

EXCITED-STATE WAVE FUNCTIONS, EXCITATION ENERGIES, AND OSCILLATOR STRENGTHS FOR Ne ($2p^5\,3s$). Albert Gold and Robert S. Knox,* Department of Physics and Institute of Optics, University of Rochester, Rochester, New York (Received October 13, 1958).

Solutions of the Hartree-Fock equations for the ^3P and ^1P terms of neon ($2p^5\,3s$) have been obtained. Wave functions are tabulated and results of computations of excitation energies and oscillator strengths are presented. The former fall within 10% of experimental values. It is found that enlarging the size of the "invariant core" used to compute the excited-state wave functions has only a small effect on the predicted energies. The predicted oscillator strength of the 736A transition is 0.11, which is in reasonable agreement with available experimental data considering the large uncertainties in the measurements. The computed diamagnetic susceptibility of the ground state is -7.4×10^{-6} cm^{-3} mole, in good agreement with experiment.

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X-RAY ABSORPTION COEFFICIENTS OF THORIUM, URANIUM, AND PLUTONIUM--EXPERIMENTAL DETERMINATION AND THEORETICAL INTERPRETATION. R. B. Roof, Jr., Los Alamos Scientific Laboratory, University of California, Los Alamos, New Mexico (Received May 26, 1958; revised manuscript received December 2, 1958).

The x-ray absorption coefficients of thorium, uranium, and plutonium have been determined for the $K\alpha$ and $K\beta$ radiations of silver, molybdenum, copper, cobalt, iron, and chromium. Four absorption edges were found in the wavelenght region covered. These were identified as the $L_{\rm I}$, $L_{\rm II}$, $L_{\rm III}$, $M_{\rm I}$ edges and a $M_{\rm II}$, $M_{\rm III}$ complex for plutonium. There is considerable disagreement between the experimental coefficients and the theoretical coefficients obtained from the Internationale Tabellen. Reasons for this discrepancy are discussed briefly. A correlation is obtained between the absorption of x-ray energy and the number of electrons in the $5f$ shell.

Through a consideration of the atomic number, a screening constant, principal quantum numbers, and the atomic energy levels for thorium, uranium, and plutonium, the x-ray absorption coefficients of these elements have been calculated. Screening constants and quantum conditions in the various equations were adjusted until, on the average, the experimental and calculated values agree within 2%. An increase in the screening constant of plutonium over that for thorium and uranium is interpreted in terms of the number of electrons occurring in the $5f$ shell. The change in certain quantum conditions for plutonium in comparison to those for thorium and uranium is interpreted in terms of plutonium being a subtransitional point within the $5f$ transition series.

X-RAY ABSORPTION COEFFICIENTS OF THE ELEMENTS WITH Z=1 TO 17 FOR Mo $K\alpha$ RADIATION. R. B. Roof, Jr., Los Alamos Scientific Laboratory, University of California, Los Alamos, New Mexico (Received July 21, 1958).

The x-ray absorption coefficients of the elements with Z = 1 to 17 (excluding helium and neon) have been determined for Mo $K\alpha$ radiation. The values found in the present work are in better agreement with the earlier experimental data of Allen than with the semitheoretical values given in the Internationale Tabellen.

EFFECT OF AN ELECTRIC DIPOLE MOMENT OF THE PROTON ON THE ENERGY LEVELS OF THE HYDROGEN ATOM. R. M. Sternheimer, Brookhaven National Laboratory, Upton, New York (Received October 8, 1958).

The perturbations of the energy levels of the hydrogen atom by a possible electric dipole moment of the proton, d, have been obtained. The inhomogeneous equation for the first-order perturbation of the wave functions was solved analytically. The shifts of the energy levels are of the second order in d. In particular, there results a decrease of the Lamb shift between the $2S_{1/2}$ and $2P_{1/2}$ levels. By equating this decrease to the maximum allowed by the approximate agreement between the experimental and theo-

retical values of the Lamb shift, it is concluded that the electric dipole moment of the proton, d, must be less than 1.3×10^{-13} cm times the charge of the electron. Calculations have also been carried out for the shifts of the energy levels of deuterium produced by a possible electric dipole moment of the deuteron.

ELECTRON SCATTERING CROSS SECTION WITH RELATIVISTIC CORRECTION BASED ON THE THOMAS-FERMI THEORY. T. Tietz, Department of Theoretical Physics, University of Łódź, Łódź, Poland (Received April 4, 1958; revised manuscript received October 21, 1958).

A formula for the total cross section for elastic scattering of fast electrons by a Thomas-Fermi potential, including relativistic effects, is derived from the differential cross section previously obtained.

X-RAY INCOHERENT SCATTERING FUNCTION OF THE ATOMS BASED ON THE THOMAS-FERMI THEORY. T. Tietz, Department of Theoretical Physics, University of Łódź, Łódź, Poland (Received October 3, 1958).

In this paper we give a simplified method of calculating the incoherent scattering function, which appears in the Thomas-Fermi theory for the x-ray. Our values for the incoherent scattering function are compared with the corresponding numerical values of Bevilogua.

FORMATION OF H^- IONS BY ELECTRON IMPACT ON H_2. G. J. Schulz, Westinghouse Research Laboratories, Beulah Road, Churchill Borough, Pittsburgh, Pennsylvania (Received August 25, 1958).

The cross section for production of H^- ions by electron impact in hydrogen gas is studied. The cross section exhibits a plateau around 10 ev with a value of 1.2×10^{-20} cm^2. A sharp peak with a cross section of 3.5×10^{-20} cm^2 is observed at 14.2±0.1 ev. The first plateau is associated with the reaction $H_2 + e \rightarrow H^- + H$, and the sharp peak with the production of hydrogen atoms in the first excited state, $H_2 + e \rightarrow H^* + H^-$.

CONTRIBUTION OF ANNIHILATION RADIATION TO THE GAMMA-RAY FLUX IN LEAD. Martin J. Berger, John H. Hubbell, and Ida H. Reingold, National Bureau of Standards, Washington, D. C. (Received October 15, 1958).

Multiply scattered gamma rays from a high-energy source, upon being absorbed, give rise to a secondary flux of annihilation radiation. A procedure is developed, within the framework of the moment method of Spencer and Fano, for calculating the production and diffusion of the annihilation radiation. The treatment includes infinite as well as semifinite media. Numerical results are presented for primary radiation sources with energies between two and ten Mev in a lead medium. These results can be represented by simple semiempirical formulas.

ELASTIC SCATTERING OF PROTONS BY NITROGEN. A. J. Ferguson, R. L. Clarke, and H. E. Gove, Atomic Energy of Canada Limited, Chalk River, Ontario, Canada (Received October 7, 1958).

The cross sections for the elastic scattering of protons by nitrogen have been measured in 105 angular distributions ranging in angle from 53° to 155° and in energy from 1.05 Mev to 2.93 Mev. Resonances have been observed at 1065 ±5 kev, 1557±6 kev, 1743±7 kev, 1803±7 kev, 2344±10 kev, and 2468±10 kev.

PHASE-SHIFT ANALYSIS OF PROTON SCATTERING BY NITROGEN. A. J. Ferguson, Atomic Energy of Canada Limited, Chalk River, Ontario, Canada (Received October 7, 1958).

A phase-shift analysis has been made of a set of angular distributions for the elastic scattering of protons by nitrogen in the energy range 1.0 Mev to 3.0 Mev. Two independent S-wave phase shifts and one P-wave phase shift of hard-sphere type have been assumed. Moderately good agreement with the nonresonant scattering below 2.3 Mev and with the scattering at the $\frac{1}{2}+$ resonance at 1.557 Mev has been obtained, indicating that the gross features of the scattering can be represented in this way. Between 2.3 Mev and 3.0 Mev the fits are poor. The results indicate a broad $\frac{1}{2}+$ resonance at 2.32 Mev with a width of 0.55 Mev.

SLOW-NEUTRON CROSS SECTIONS OF Pu^{240}, Pu^{242} AND Am^{243}. R. E. Coté, L. M. Bollinger, R. F. Barnes, and H. Diamond, Argonne National Laboratory, Lemont, Illinois (Received October 9, 1958).

Measurements of the total cross sections of Pu^{240}, Pu^{242} and Am^{243} have been made using the Argonne "fast chopper." The results on Pu^{240} include only the resonance structure which could be studied through the Pu^{240} impurity in some samples of Pu^{239} which were used in an exhaustive study of the latter nuclide. The values of the resonance parameters of the resonances at 1.053, 20.4, and 38.2 ev comprise these results. The studies of Pu^{242} and Am^{243} were made possible by the separation of these nuclides from samples of Pu^{239} which had undergone extensive neutron irradiation. Resonances were observed at 2.65 and 53.6 ev in Pu^{242} and at 0.976, 1.353, 1.74, 3.42, 5.12, 6.54, 7.84, 10.3, 12.8, 13.1, and 15.3 ev in Am^{243}. Resonance parameters are reported for all of these resonances and, in addition, the resonance capture integrals derived from these parameters are compared with the results of other measurements of these quantities.

GAMMA-RAY EXCITATION OF THE 15.1-Mev LEVEL IN C^{12}. E. L. Garwin, The Enrico Fermi Institute for Nuclear Studies, The University of Chicago, Chicago, Illinois (Received October 13, 1958).

The parameters of the 15.1-Mev level in C^{12} have been determined. The angular distribution of scattered photons has been determined to be dipole. A scatterer-associated background to the scattering from this level has been demonstrated. The peak absorption cross section of 29.7 ± 1.1 barns has been determined from a self-absorption experiment. The self-absorption measurement and a production experiment combine to give values of 64.5 ± 10.4 ev and 59.2 ± 9.7 ev for the total level width and radiative width to the ground state, respectively. The combination of the value of the integrated elastic scattering cross section from the absolute measurement, and those values calculated from the self-absorption and production experiments, gives for this integrated cross section the value 2.33 ± 0.19 mb-Mev. The radiation width to the 4.43-Mev state of C^{12} is determined to be 3.2 ± 2.5 ev by fitting the detector resolution function to the observed pulse-height spectra. A value for the inelastic scattering in the giant resonance region of C^{12} is determined.

STUDY OF LEVELS IN O^{18} THROUGH THE RADIATIVE CAPTURE OF ALPHA PARTICLES BY C^{14}. H. E. Gove and A. E. Litherland, Chalk River Laboratory, Atomic Energy of Canada Limited, Chalk River, Ontario, Canada (Received October 10, 1958).

Two resonances in the reaction $C^{14}(\alpha,\gamma)O^{18}$ have been studied at laboratory alpha-particle energies of 1.142 ± 0.010 and 1.790 ± 0.010 Mev. These correspond to excited states in O^{18} at 7.127 and 7.630 Mev. At the upper resonance only two primary transitions are observed leading to the ground state of O^{18} and the first excited state at 1.98 Mev. Their angular distributions along with other evidence unambiguously establish spins and parities of 1- for the 7.63-Mev state, 2+ for the 1.98-Mev state, and 0+ for the ground state. Values of $\gamma = \Gamma_\alpha \Gamma_\gamma / \Gamma$ of 80 and 160 millielectron volts, respectively, are obtained for the 7.63- and 5.65-Mev $E1$ primary transitions. At the lower resonance again only two primary transitions are observed leading to levels in O^{18} at 1.98 and 3.55 Mev. Greater than 96% of the decays of the 3.55-Mev level lead to the 1.98-Mev level. Analysis of the angular distributions of the four gamma rays observed in the direct spectrum at this resonance, with respect to the incident beam, unambiguously establish spin and parity of 4+ to both the capturing state at 7.13 Mev and the level at 3.55 Mev. Values of γ of 15 and 12 millielectron volts respectively are obtained for the 3.58-Mev $M1$ and the 5.15-Mev $E2$ primary transitions. The amount of $E2$ mixing in the former transition is very small.

RADIATIONS FROM Ba^{129} Wolfgang Henkes, Chemistry Department, Brookhaven National Laboratory, Upton, New York (Received October 8, 1958).

The decay of Ba^{129} produced by an $(n, 2n)$ reaction has been investigated. A new half-life value of 2.45 ± 0.05 hours is reported. Gamma rays of 1.45, 0.210, 0.182, and 0.127 Mev with relative intensities 42, 65, 100, and ~26 were found, with coincidences between the 1.45-Mev and 0.182-Mev gammas. Milking experiments showed 25% to 30% of Ba^{129} to decay by β^+ emission,

based on assumptions on the unknown decay scheme of Cs^{129}. This result is compatible with the intensities of K x-rays, positrons, and the various gamma rays only if unobserved γ-γ coincidences are postulated in addition to the reported ones. The result of the milking experiments is also in contradiction to the calculated K/β^+ ratio. 19% of the observed β^+ transitions are in coincidence with the 0.210- and 5.6% with the 0.127-Mev gamma rays. Because of the mentioned discrepancies, no decay scheme can be proposed.

RELATIVE PULSE HEIGHT OF PROTONS AND ELECTRONS IN KI(Tl). Paul Kienle, * Brookhaven National Laboratory, Upton, New York, and R. E. Segel, † Aeronautical Research Laboratory, Dayton, Ohio (Received July 23, 1958).

The relative scintillation efficiency for protons and electrons in KI(Tl) has been measured in the energy range of a few Mev. Protons were found to give a greater light output than electrons, and a pulse-height ratio $p{:}e = 1.42{:}1$ was determined.

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DIRECTIONAL CORRELATION OF GAMMA-RAYS FOLLOWING THE DECAY OF Eu^{152}. R. W. Lide, * and M. L. Wiedenbeck, Department of Physics, The University of Michigan, Ann Arbor, Michigan (Received October 2, 1958).

Directional correlation measurements have been made on the 245-kev—122-kev, 969-kev—122-kev, 1118-kev—122-kev, 1416-kev—122-kev, 872-kev—245-kev, and 1170-kev—245-kev cascades in Sm^{152}, and on the 782-kev—345-kev cascade in Gd^{152}. The spins of Sm^{152} are found to be 0, 2, 4, 2, 3, 3 for the ground, 122-kev, 367-kev, 1092-kev, 1240-kev, and 1538-kev levels. The 245-kev gamma ray is found to be pure quadrupole; the 872-kev gamma ray is 98% quadrupole, 2% dipole; the 969-kev gamma ray is 98% quadrupole, 2% dipole; the 1118-kev gamma ray is 99.8% quadrupole, 0.2% dipole; the 1170-kev gamma ray is 98% quadrupole, 2% dipole; the 1416-kev gamma ray is 15% quadrupole, 85% dipole. In Gd^{152} the spin of the 1127-kev level is found to be 3 and the 782-kev gamma ray is found to be pure dipole. A spin assignment of 1, 2, or 3 for the 757-kev level would not be inconsistent with the data. For any of these cases the 412-kev gamma ray must be mostly dipole; the quadrupole content must be less than 15%.

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DECAY OF Pm^{143} AND Pm^{144}. Shimon Ofer, * Brookhaven National Laboratory, Upton, New York (Received October 14, 1958).

Radioactive Pm^{143} and Pm^{144} were produced by the bombardment of Pr^{141} with α particles. Gamma rays at 475 kev, 610 kev, and 695 kev are found to be associated with the decay of Pm^{144} by K-capture to Nd^{144}. A 740-kev gamma transition is associated with the decay of Pm^{143} by K-capture to Nd^{143}. Levels of Nd^{144} are placed at 695 kev (2+), 1305 kev (4+), and 1780 kev (6+) above the ground level. Assignments of spins of levels and multipolarities of gamma-ray transitions are based on the results of measurements of conversion coefficients and gamma-gamma angular correlations. A value of $(13\pm4)\times10^{-11}$ sec was found for the mean life of the 1305-kev level of Nd^{144}. An upper limit of 3×10^{-11} sec was found for the mean life of the 695-kev level of Nd^{144}.

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NEUTRON CAPTURE GAMMA RAYS IN Cl^{36}. R. E. Segel, * Aeronautical Research Laboratory, Wright-Patterson Air Force Base, Dayton, Ohio (Received October 13, 1958).

γ-γ coincidences on the gamma rays following thermal neutron capture in Cl^{35} have been measured. Combining these results with the energy levels in Cl^{36} known from the $Cl^{35}(d,p)Cl^{36}$ reaction and the $Cl^{35}(n,\gamma)Cl^{36}$ gamma-ray spectrum measured by other workers, a decay scheme is constructed which unambiguously places most of the known gamma rays in Cl^{36}. An examination is made of the reduced widths of gamma rays emanating from the capturing state, and it is shown that the reduced widths for gamma rays of the same multipolarity can fluctuate widely, and that these fluctuations do not appear to be correlated with the final-state shell model configuration. It is also shown that the reduced widths for $E1$ and $M1$ transitions emanating from the capturing state are both significantly smaller

than that calculated from the single-particle estimate, and that $E1$ transitions are the more intense by about a factor of four. Evidence is presented for there being collective motion present in some of the higher excited states in Cl^{36}.

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PHOTONEUTRON CROSS SECTIONS OF Li^6 AND Li^7. T. A. Romanowski* and V. H. Voelker, Case Institute of Technology, Cleveland, Ohio (Received April 24, 1958; revised manuscript received October 17, 1958).

The total photoneutron cross sections for Li^6 and Li^7 were determined from the yields measured with a neutron detector. The efficiency of the neutron detector was found from the comparison of the detected and computed neutron yield from copper. In the calculation of the cross section, the inverse photon matrix formulation of the photon difference method was used. The Li^6 cross section has a peak at about 12.5 Mev and its absolute value at the peak is equal to 2.6 ± 0.53 mb. The Li^7 cross section has a peak at about 14 Mev and its value at the peak is equal to 3 ± 0.75 mb.

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COULOMB CORRECTIONS IN THE THEORY OF INTERNAL BREMSSTRAHLUNG. Larry Spruch and Wallace Gold,* Physics Department, New York University, Washington Square, New York, New York (Received October 7, 1958).

The internal bremsstrahlung associated with allowed β decay is calculated for the case for which the gamma-ray energy is less than $2mc^2$, the kinetic energy of the electron in its final state is small compared to mc^2, and $Ze^2/\hbar c$ is small compared to 1. It is not assumed that $(Ze^2/\hbar c)mc$ is small compared to the final momentum of the electron, or to the gamma-ray momentum. Results are obtained for the gamma-ray energy spectrum and for the angular correlation between the electrons and the gamma rays. For S^{35}, for which the above assumption would seem to be satisfied, the agreement between theory and experiment for the number of gamma rays per β disintegration per mc^2 is better than that previously obtained; due to uncertainties in the experimental results, the extent of the improvement is not clear. Under the additional assumption that the final kinetic energy of the electron is small compared to the gamma-ray energy, an expression is derived for the polarization of the gamma rays.

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SLOW-NEUTRON SCATTERING BY ROTATORS. Howard C. Volkin, Lewis Flight Propulsion Laboratory, National Advisory Committee for Aeronautics, Cleveland, Ohio (Received October 10, 1958).

The methods of Zemach and Glauber for treating the scattering of low-energy neutrons by molecules are extended for the rotational degrees of freedom to any type of rotator. A procedure is developed which enables their results to be generalized to symmetric or asymmetric molecules. The method is then applied to the calculation of the direct scattering up to the accuracy of the first quantum-mechanical correction to the classical cross section.

CHARGE DISTRIBUTION OF EXCITED ISOMERIC NUCLEI AND ATOMIC SPECTRA (THE NUCLEAR ISOMERIC SHIFT). Richard Weiner, Physical Institute of the Academy, Bucharest, Rumania (Received March 31, 1958; revised manuscript received September 22, 1958).

Attention is called to an electric characteristic of excited nuclei almost unstudied until the present, neither theoretically nor experimentally – namely the charge distribution. The comparison of the charge distributions of two isomeric nuclei is automatically realized in atomic spectra giving rise to the nuclear isomeric shift of spectral lines.

Some general theoretical aspects of this effect are here discussed. Only odd nuclei with optical protons are studied. It is assumed that (A) the transitions are single-particle ones; (B) the Rosenthal-Breit perturbation theory is valid.

Under these assumptions it is shown that but for the sign, the effect is a pure single-particle effect, given by the optical protons. The sign of the shift is generally intimately related to the whole nuclear configuration. In the case of two most characteristic transitions it is shown that the order of magnitude of the effect

does not depend on the shape of nuclear potential. By specializing for two particular forms of nuclear potential (harmonic oscillator and infinite square well), it is shown that there exists a very simple relation between the shift and the characteristics of the two nuclear states involved. Numerical applications are given for $In^{115}III$ and $Au^{197}I$. The theoretical value of the effect is surely within the reach of atomic spectroscopy ($\geq 10^{-2}\,cm^{-1}$).

DECAY OF Nb^{92}. Harry I. West, Jr., Lloyd G. Mann, and Glen M. Iddings, University of California Radiation Laboratory, Livermore, California (Received March 26, 1958).

The decay scheme of Nb^{92} has been reexamined. Levels at 0.934 Mev (97.4% per disintegration) and 1.82 Mev (2.6%) were found. The 1.82-Mev state decays by 32% to the ground state and 68% to the 0.934-Mev state. This is in agreement with the work of Hayward, Hoppes, and Ernest. In addition, a weak positron branch of $(5.6\pm0.6)\times10^{-2}$% per disintegration was found going to the 0.934-Mev level. Also, the angular correlation of the gamma-ray cascade was measured and found to be uniquely consistent with a level assignment of 2^+, 2^+, 0^+, with essentially pure $M1$ radiation ($E2$ admixture <0.2) for the 2^+, 2^+ transition. The large amount of $M1$ radiation in this type of transition is quite exceptional. Simple theoretical considerations indicate the gamma-ray transitions to be of the single-particle type.

REACTION $A^{40}(p,n)K^{40}$ AND THE DECAY OF K^{40}. R. E. Holland and F. J. Lynch, Argonne National Laboratory, Lemont, Illinois (Received August 11, 1958).

The energy difference involved in the decay of K^{40} to A^{40} by electron capture was measured in two ways. First, a time-of-flight technique was used to observe the neutron spectra from $A^{40}(p,n)K^{40}$ and the corresponding Q's were computed. Second, the thresholds for production of certain gamma rays from this reaction were measured; these thresholds provided another set of values for the Q's. When these Q values were combined with our measurements of the gamma-ray energies, we obtained a mass difference equivalent to 1.522 ± 0.006 Mev between A^{40} and K^{40} and an energy release of 60 ± 8 kev in the decay of K^{40} to A^{40}. This appears to conflict with other information on this branch of the decay of K^{40}.

25-MINUTE ISOMER OF Se^{83}. R. G. Cochran and W. W. Pratt, The Pennsylvania State University, University Park, Pennsylvania (Received July 14, 1958).

Se^{83} was produced by irradiation of a selenium sample, enriched to 75% in the Se^{82} isotope, in the Pennsylvania State University Research Reactor. The gamma-ray spectrum of the 25-minute isomer was investigated by means of a 3 in. $\times$3 in. diameter NaI(Tl) crystal scintillation spectrometer. Gamma rays of energy 2.294 $\pm$0.030 Mev, 1.880$\pm$0.015 Mev, 1.309$\pm$0.005 Mev, 1.058$\pm$0.005 Mev, 0.833$\pm$0.005 Mev, 0.712$\pm$0.010 Mev, 0.524$\pm$0.012 Mev, 0.358$\pm$0.005 Mev, and 0.225$\pm$0.005 Mev were found. The beta-ray spectrum was investigated by means of a plastic scintillator. In order to eliminate a competing activity from Se^{81} it was found essential to restrict the observation of beta rays to those in coincidence with the Se^{83} gamma-ray spectrum. A coincidence circuit of 0.25-microsecond resolving time was developed for this purpose. The beta-ray spectrum was found to be complex with end-point energies of approximately 0.45 Mev, 1.0 Mev, and 1.7 Mev. From these data, together with gamma-gamma coincidence measurements, a decay scheme is proposed.

$C^{12}(n,p)B^{12}$ CROSS SECTION FOR 14.9- TO 17.5-Mev NEUTRONS. William E. Kreger and Bernard D. Kern, * U. S. Naval Radiological Defense Laboratory, San Francisco, California (Received September 2, 1958).

The cross section for the $C^{12}(n,p)B^{12}$ reaction has been measured for 14.9- to 17.5-Mev neutrons. The neutrons were obtained from the $T(d,n)He^4$ reaction and their flux density was determined by counting the recoil alpha particles or by counting the neutrons directly with a $Li^6I(Eu)$ scintillation counter. A cylindrical plastic scintillator 5 inches in diameter and 3 inches in length served as the carbon-containing target and permitted the counting of the B^{12} decay beta rays in nine consecutive 7-msec intervals during the "beam off" period of a pulsed neutron beam cycle. The cross section rises from slightly above the reaction threshold of 13.6 Mev to a value of 29.09 $\pm$4.36 millibarns at 17.5 Mev. The B^{12} beta-decay

half-life has been redetermined as 18.87±0.50 milliseconds.

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RADIATIONS FROM Tb^{156}. J. T. Holloway* and L. Jackson Laslett, Institute for Atomic Research and Department of Physics, Iowa State College, Ames, Iowa (Received August 4, 1958).

The decay of Tb^{156} has been studied following its production by the $Tb^{159}(\gamma, 3n)Tb^{156}$ reaction with 64-Mev bremsstrahlung. In the 5.6-day activity, 18 gamma-transitions were detected with energies between 89 kev and 2060 kev. The K x-radiation of Gd was associated with this activity, while no Dy radiation was found. No annihilation radiation was detected. The gamma spectrum was observed in coincidence with the Gd x-ray and found to be closely similar to that observed alone. No evidence of β-decay was found and it is concluded that the observed decay is by orbital electron capture. A relatively weak 5.5-hour activity, ascribable to Tb^{156m}, was also produced by the $(\gamma, 3n)$ reaction.

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COINCIDENCE STUDIES IN THE DECAY OF Tb^{156} AND Tb^{154}. R. W. Henry, L. T. Dillman, N. B. Gove, and R. A. Becker, Physics Department, University of Illinois, Urbana, Illinois (Received August 15, 1958).

Tb^{156} (5.35 days) and Tb^{154} (7.5 hours and 21.5 hours) were produced by $(\gamma, 3n)$ and $(\gamma, 5n)$ reactions on stable Tb^{159}. In addition, Tb^{156} was produced by a $(p, 2n)$ reaction on stable Gd^{157}. Gamma rays at 89±1, 200±2, 261±5, 300±10, 358±5, 420±5, 535±3, 1050±10, 1140±10, 1210±10, 1410 ±10, 1630±20, and 1830±20 kev were associated with the 5.35-day activity, and 88±1, 123±3, 180 ±5, 250±3, 340±10, 511±10, 640±15 and others greater than 1600 kev were associated with 21.5-hour and shorter activities. Levels of Gd^{156} are placed at 89 (2+), 289 (4+), 1140, 1500, and 2035 kev above the ground state. These follow electron capture in 5.35-day Tb^{156}. The K-shell conversion coefficients of the 89- and 200-kev transitions are calculated to be about 1.0 and 0.16, respectively.

ANTIPROTON-NUCLEON ANNIHILATION PROCESS. II. Owen Chamberlain, Gerson Goldhaber, Louis Jauneau, Theodore Kalogeropoulos, Emilio Segrè, and Rein Silberberg, Radiation Laboratory and Department of Physics, University of California, Berkeley, California (Received October 13, 1958).

We have continued work on antiproton interactions in photographic emulsions. Most of the data come from an exposure at the Bevatron to an enriched antiproton beam of momentum 700 Mev/c. In this paper we present the analysis of 221 antiproton stars, 95 of which occurred in flight. We find an average antiproton cross section of $(1.9\pm 0.2)\sigma_0$, where $\sigma_0 = \pi(1.2\times 10^{-13}A^{1/3})^2$ cm^2, for all the elements in emulsion, excluding hydrogen. The primary antiproton annihilation gives rise to 5.36 ± 0.3 pions on the average. In annihilations at rest, 1.3 of the pions formed interact with the parent nucleus; in reactions in flight, 1.9 of the pions. For stars at rest, the energy available in the annihilation in complex nuclei is divided up among the products as follows: charged pions, 48± 6%; neutral particles (other than neutrons and K^0 mesons) 28 ± 7%; K mesons 3 ± 1.5%; and cascade nucleons and nuclear excitation 21± 2%. For the stars in flight the corresponding percentages are 45 ± 7%, 22 ± 7%, 3 ± 1.5%, and 30± 2%. To fit the average pion multiplicity, the interaction radius of the Fermi statistical model must be taken as $2.5\hbar/m_\pi c$. Other proposals to explain the large multiplicity are discussed. We deduce from the fraction of pions interacting in the same nucleus that the annihilation takes place at the outer fringes of the nucleus.

EMISSION OF HEAVY FRAGMENTS IN NUCLEAR DISINTEGRATIONS. O. Skjeggestad and S. O. Sörensen, Institute of Physics, University of Oslo, Blindern, Norway (Received October 7, 1958).

A phenomenological description is given of the processes which lead to the emission of energetic heavy fragments from cosmic-ray stars. The energy spectra, angular distribution, and relative frequency of the particles have been determined and a detailed comparison is made with the predictions from nuclear evaporation theory. A simple method for charge determination of the fragments is described.

DECAY OF HYPERONS AND MESONS FROM THE UNIVERSAL FERMI INTERACTION. Akihiko Fujii and Masaaki Kawaguchi, Department of Physics, Purdue University, Lafayette, Indiana (Received October 15, 1958).

The decay of the hyperon, charged pion, and K meson (except $K_{\pi 3}^{+}$) is investigated on the basis of the universal V-A Fermi interaction together with the idea of the Gell-Mann tetrahedron, by treating the virtual baryon-antibaryon pair effect in a phenomenological way. It is shown that the decay rate of $K_{\pi 2}^{+}$ calculated by the parameters adjusted to the decay of π^{+}, $K_{\mu 3}^{+}$, and K_{e3}^{+} is in agreement with experiment, thus suggesting a possible consistent picture of the model.

COVARIANT CONSERVATION LAWS IN GENERAL RELATIVITY. Arthur Komar, Physics Department, Syracuse University, Syracuse, New York (Received October 15, 1958).

A set of covariant conservation laws is constructed in the general theory of relativity. Their relationship to the generators of infinitesimal coordinate transformations is indicated. In a given coordinate system certain of these quantities may be naturally identified as energy and momentum. We can continue to recognize these conserved quantities in all coordinate systems because of the covariant character of the expressions.

V-A THEORY AND THE DECAY OF THE Λ HYPERON. S. Okubo and R. E. Marshak, University of Rochester, Rochester, New York, and E. C. G. Sudarshan, University of Rochester, Rochester, New York, and Harvard University, Cambridge, Massachusetts (Received September 8, 1958).

The decay of the Λ hyperon is studied within the framework of the chirality-invariant four-fermion interaction. It is shown that the branching ratio of the charged and neutral modes, the s- to p-wave emission, as well as the sign and magnitude of the asymmetry parameter of the $p+\pi^{-}$ decay mode can be understood on the basis of the V-A theory. Improvements upon the Born approximation, using dispersion theory, indicate that these conclusions are not invalidated by taking into account the pion-nucleon interaction.

GENERAL TRANSFORMATION OF THE SYMMETRICAL PSEUDOSCALAR THEORY. A. Pais, Institute for Advanced Study, Princeton, New Jersey, and R. Serber, Department of Physics, Columbia University, New York, New York (Received October 1, 1958).

The symmetrical pseudoscalar theory with extended source is subjected to a rigorous transformation in such a way that the transformed Hamiltonian depends explicitly on the total isotopic spin and total angular momentum of the system. These collective variables are introduced by the same method of employing a general distribution function which was previously studied by the authors for the charged scalar theory.

THEORIES OF Σ^{-}-p CAPTURE. Marc Ross, Physics Department, Indiana University, Bloomington, Indiana (Received October 1, 1958).

Capture rates for production of Σ^{0} and Λ^{0} from the zero-energy Σ^{-}-p system are calculated in Born approximation on the basis of two theories of hyperon-nucleon forces. In one theory, the relative (Λ^{0}, Σ) parity is even, and pion-baryon interactions are symmetric. In the other theory the (Λ, Σ) parity is odd and there is a linear scalar π-Λ^{0}-Σ interaction. According to the results, the capture takes place from the $2P$ orbit and the calculated Λ^{0}/Σ^{0} ratio is compatible with the experimental ratio of two in both the even- and odd-parity cases.